ABSTRACT

Title of dissertation: NEW MECHANISMS FOR TRANSMISSION OF SUPERSYMMETRY BREAKING

Siew Phang Ng, Doctor of Philosophy, 2004

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We considered new mechanisms for transmission of supersymmetry breaking and their phenomenological consequences. Specifically, we investigated the scalar mass corrections via five-dimensional supergravity loops and explored the possibility that supersymmetry is an accidental symmetry of Nature. We find that both these lead to phenomenologically viable scenarios. In the former, the negative slepton mass-squared masses arising from minimal anomaly mediation is cured. In the latter, which constitutes a paradigm shift, we find a more natural framework for low-energy supersymmetry than the conventional picture.

NEW MECHANISMS FOR TRANSMISSION
OF SUPERSYMMETRY BREAKING

by

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DEDICATION

To my parents.

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Chapter 1: Introduction

In the 1960's particle physicists were confronted with a menagerie of particles and no coherent unified framework to understand the subatomic world. With significant theoretical developments [1] in the latter part of the decade and early in the next, it soon became apparent that the subatomic phenomena in our universe could be described in terms of gauge theories. These gauge theories possess a built-in redundancy in their description of physical phenomena in such a way that it is possible to choose any description (historically termed a gauge) and get the same result when one calculates any physical observable regardless of the gauge choice. Underlying this powerful result is the concept of the gauge symmetries. These gauge symmetries, however, constrain the types of interactions one can have in a theory. In particular, explicit mass terms for quarks are forbidden. There are a few ways to give mass to a quark but arguably, the most elegant method is to postulate the existence of a scalar (spin-0) particle called the Higgs. What the Higgs does is that it couples to two quarks and assuming electroweak symmetry breaking, the Higgs particle acquires a vacuum expectation value (vev) thereby dynamically generating a mass term for the quark. There are, however, inherent dangers in this approach as the Higgs particle, unprotected by any symmetry, can acquire a heavy mass through radiative corrections. The fact that electroweak symmetry breaking occurs at $\mathcal{O}(100)$ GeV while the Planck scale is at $\mathcal{O}(10^{18})$ GeV leads one to conclude that in the absence of additional symmetries or clever mechanisms, there is an extreme degree of fine-tuning commonly known as "The Hierarchy Problem".

This is where supersymmetry [2] comes in. Supersymmetry (or SUSY as it is affectionately known to its practitioners) requires that for every type of particle that we have seen, there exists a corresponding supersymmetric partner with exactly the same properties except it would have a different spin. The fermionic quark would have a supersymmetric partner called squark that is bosonic while the supersymmetric partner of a bosonic photon would be a fermionic photino. When SUSY was discovered in the early 1970's , it did not take particle physicists long to realize that with a supersymmetric partner for every known ordinary particle, the radiative corrections coming from the ordinary particles would be nearly completely cancelled by the radiative corrections coming from its supersymmetric analogs. The degree of cancellation would then depend on how badly broken SUSY is. One might wonder why we need to break SUSY. The fact of the matter is that we simply do not observe any of the supersymmetric partners.

The simplest possibility is that supersymmetry is broken and therefore the supersymmetric partners can have properties that differ from that of their ordinary counterparts. As we have experimentally probed only low energies, we would expect that supersymmetric partners, if present, to have masses higher than their ordinary analogs. There are a variety of ways to break supersymmetry but these do not usually provide us with definitive predictions as the mode of transmission of the supersymmetry breaking can significantly alter the result. The point of view we adopt is that we assume that supersymmetry is broken in some hidden sector and we look at how this is transmitted to the visible sector we live in. A very compelling idea that perfectly complements the existence of SUSY is the brane-world scenario

[3] where it is postulated that the universe is higher-dimensional and contains in addition to the ordinary single temporal and three spatial dimensions, extra spatial dimensions that are are compactified and orbifolded. Essentially, the orbifold turns, say, a circle into an interval by "modding out" half the circle. We identify the endpoints of the interval as branes, where we take one to represent our visible sector and the other as the hidden one, and the space in between the branes as the bulk. There are, of course, many different ways of transmitting SUSY breaking across the extra dimension but in this thesis, we shall investigate the transmission via five-dimensional supergravity loops. A few words are in order here. Supergravity [4] is the gauge theory of the supersymmetry algebra which, because it contains a spin-2 particle, is identified as a theory of gravitons and necessarily gravitinos (the spin-$\frac{3}{2}$ supersymmetric partner of the gravitons). An interesting property of supergravity is that it is ubiquitous and permeates the entire bulk because as long as we have space-time, we have fluctuations of space-time, *i.e.* gravitons. Hence, there is considerable interest in working out the radiative corrections of this setup. In Chapter 2, we will consider these effects in a standard brane-world scenario while in Chapter 3, we deal with a variety of more exotic scenarios. A new formalism for calculating these loop effects is also presented.

Another possible resolution to the fact that we have yet to see supersymmetric particles lies in the possibility that supersymmetry is not a fundamental symmetry of nature but rather an accidental symmetry much like baryon number. The big picture is that our universe lies in some superconformal basin with an attractive fixed point in the middle. As the energy is evolved downwards, degrees of freedom get integrated out because they become extremely heavy with respect to the processes at that energy and hence decouple from the theory. This would correspond to our universe rolling down from the edge of the superconformal basin towards the low energy fixed point. Hence, the theory would become more and more supersymmetric as we go down in energy thereby giving us the same physics as the case where SUSY is an exact symmetry at high energies but gets broken as one goes down in energy. However, we do not currently possess the theoretical understanding to analyze such a strongly coupled superconformal field theory. Instead, we make use of the Anti de-Sitter-Conformal Field Theory (AdS-CFT) correspondence [5] which relates the picture we have just painted with an extra-dimensional Anti de-Sitter setup. By exploiting the correspondence, we can construct these theories on the AdS side where it is relatively straightforward to get a handle on things. The construction and the phenomenological consequences of the resulting model are explored in Chapter 4.

Chapter 2: Supergravity Loop Contributions to Brane World Supersymmetry Breaking

We compute the supergravity loop contributions to the visible sector scalar masses in the simplest 5D 'brane-world' model. Supersymmetry is assumed to be broken away from the visible brane and the contributions are UV finite due to 5D locality. We perform the calculation with $\mathcal{N} = 1$ supergraphs, using a formulation of 5D supergravity in terms of $\mathcal{N} = 1$ superfields. We compute contributions to the 4D effective action that determine the visible scalar masses, and we find that the mass-squared terms are negative.

2.1 Background

In this chapter, we study supersymmetry (SUSY) breaking in the simplest 5D 'brane world' scenario. In brane world scenarios, some or all of the visible sector fields are assumed to be localized on a brane, and SUSY is broken away from the visible brane. In this case, bulk fields transmit the message of SUSY breaking to the visible sector. We consider the minimal case where the bulk fields are the 5D supergravity (SUGRA) multiplet. Thus, supergravity plays the role of the messenger for SUSY breaking. Previously, Ref. [6] showed that the leading contribution to visible sector SUSY breaking for large radius, comes from anomaly-mediated SUSY breaking (see also Ref. [7]). If the visible sector consists only of the minimal supersymmetric standard model, the slepton mass-squared terms are negative. Thus, for these brane-world models to be realistic we require other contributions to SUSY breaking in the visible sector. With the hope of getting positive mass-squared terms, we will calculate the leading contributions to SUSY breaking by SUGRA loops.

The simplest 5D brane-world scenario can be described as follows. The 5D space-time is flat and compactified on an S^1/Z_2 orbifold. There is one 3-brane at each of the Z_2 fixed points. These 3-branes can be regarded as the boundaries of the extra dimension of length $\ell = \pi r$, where r is the radius of the S^1. We assume that SUSY is broken by the vacuum expectation value of a chiral superfield X localized on the hidden brane. The visible chiral superfields Q are assumed to be localized on the other brane. In this 5D effective theory, contact terms between Q and X are forbidden by 5D locality.[1] The effects of SUGRA mediated SUSY breaking can be analyzed systematically using the 4D effective Lagrangian that describes the physics below the compactification scale $1/r$. The effective theory contains the chiral superfields Q and X, the 4D SUGRA multiplet, and the chiral radion multiplet

$$T = \pi r + \cdots + \theta^2 F_T. \tag{2.1}$$

Expanding the 4D effective action in Q and X, the leading terms involving Q that cannot be forbidden by symmetries are

$$\Delta\mathcal{L}_{4,\text{eff}} = \int d^4\theta \left[c_1(T) Q^\dagger Q + c_2(T) X^\dagger X Q^\dagger Q + \cdots \right]. \tag{2.2}$$

At tree level, c_1 is independent of T, and Ref. [8] showed that c_2 vanishes. Therefore, we must consider loop effects.

[1] In a more fundamental theory with additional states with masses $M \gg 1/r$, contact terms between Q and X will be suppressed by e^{-Mr}.

At 1-loop level, there are contributions to c_1 from the diagrams in Fig. 1. These contributions are of order

$$c_1 \sim \frac{1}{M_5^3(T+T^\dagger)^3}. \tag{2.3}$$

The dependence on T is fixed by dimensional analysis and the observation that c_1 cannot depend on the fifth component of the graviphoton of 5D SUGRA, which is contained in $T - T^\dagger$ [8]. Loop corrections to c_1 are finite because they are sensitive to the size of the extra dimension, while all divergent effects are local. These corrections can give important contributions to the scalar masses of Q if $\langle F_T \rangle \neq 0$:

$$\Delta m_Q^2 = -3\langle c_1 \rangle \left| \left\langle \frac{F_T}{T} \right\rangle \right|^2. \tag{2.4}$$

Here we have neglected 1-loop operators of the form

$$\Delta\mathcal{L}_{4,\text{eff}} \sim \int d^4\theta \, \frac{|D^2 T|^2}{M_5^3(T+T^\dagger)^2} Q^\dagger Q, \tag{2.5}$$

which give contributions to the scalar masses proportional to $\langle F_T \rangle^4$. Thus, the contribution from c_1 in Eq. (2.4) dominates only if $\langle F_T \rangle \ll 1$. A nonzero value for $\langle F_T \rangle$ is equivalent to the Scherk–Schwarz [9] mechanism for SUSY breaking, as discussed in Ref. [10]. The SUGRA loop effect proportional to c_1 was computed in Ref. [11] using the off-shell formulation of supergravity due to Zucker [12]. It was found that the resulting scalar mass-squared terms are negative.

There are 1-loop contributions to c_2 from the diagrams in Fig 2. These diagrams are UV finite because the loop cannot shrink to zero size. By dimensional analysis, these give

$$c_2 \sim \frac{1}{M_5^6(T+T^\dagger)^4}. \tag{2.6}$$

This is suppressed by extra powers of M_5 compared to c_1. This contribution may be important if $\langle F_T \rangle$ is sufficiently small. In this case, it gives a contribution to the Q scalar mass

$$\Delta m_Q^2 = -\langle c_2 \rangle |\langle F_X \rangle|^2. \tag{2.7}$$

Although c_1 is known in the literature, c_2 has never been calculated. In this chapter, we will present explicit calculations of both c_1 and c_2. We perform quantum computations using supergraphs (see e.g. [13], [14]) applied to the formulation of 5D SUGRA in $\mathcal{N} = 1$ superspace developed in Ref. [15]. This formalism has several advantages over component calculations. First, higher powers of Dirac delta functions from brane-bulk interactions do not arise in this formulation. Higher powers of Dirac delta functions occur only after integrating out auxiliary fields [43], and therefore are absent in supergraph calculations. Furthermore, the gauge can be fixed so that the superspace supergravity propagator has the following trivial tensor structure:

$$\langle V_m V_n \rangle \sim \frac{\eta_{mn}}{\Box_5}, \tag{2.8}$$

where $m, n = 0, \ldots 3$ are 4D Lorentz indices and V_m is the SUGRA superfield prepotential. The simple form of this propagator makes quantum calculations straightforward. Another advantage of this approach is that we only need to calculate five super Feynman graphs. In a direct component formulation, this number would grow by an order of magnitude.

This chapter is organized as follows. Section 2 reviews 5D SUGRA in $\mathcal{N} = 1$ superspace [15], and proves the existence of the remarkably simple gauge fixing noted above.

Section 3 gives the supergraph Feynman rules for the theory. Sections 4 and 5 contain the calculations of c_1 and c_2, respectively. We find that both c_1 and c_2 give negative scalar mass-squared terms in the visible sector. The result for c_1 agrees with Ref. [11], while the result for c_2 is new.

2.2 Lagrangian and Gauge-fixing

The Lagrangian for linearized minimal 5D SUGRA was written in terms of $\mathcal{N} = 1$ superfields in Ref. [15]. Here, we describe the component field embedding and state the superfield action. We then prove the existence of the gauge choice Eq. (2.8).

The formulation of Ref. [15] contains two real superfields V_m and P, a chiral superfield T, and an unconstrained superfield Ψ_α.[2]

The embedding of the 5D propagating fields into these superfields is accomplished as follows. The graviton, graviphoton and gravitino are first dimensionally reduced:

$$\begin{aligned} &h_{MN} \to h_{mn},\ h_{5m}, h_{55}, \\ &B_M \to B_m, B_5, \\ &\psi_{M\tilde{\alpha}} \to \psi^{(\pm)}_{m\alpha}, \end{aligned} \tag{2.9}$$

Here the 5D gravitino is decomposed into components with parity ± 1 under the Z_2

[2]The field Ψ_α corresponds to what was called $\hat{\Psi}_\alpha$ in Ref. [15].

transformation $x^5 \mapsto -x^5$. These reduced fields are embedded in superfields as

$$
\begin{aligned}
V_m &\sim \theta\sigma^n\bar{\theta}h_{mn} + \bar{\theta}^2\theta^\alpha\psi^{(+)}_{m\alpha} + \cdots, \\
\Psi_\alpha &\sim \bar{\theta}^{\dot{\alpha}}(B_{\alpha\dot{\alpha}} + ih_{5\alpha\dot{\alpha}}) + \theta\sigma^m\bar{\theta}\psi^{(-)}_{m\alpha} + \bar{\theta}^2\psi^{(-)}_{5\alpha} + \cdots, \\
T &\sim h_{55} + iB_5 + \theta^\alpha\psi^{(+)}_{5\alpha} + \cdots.
\end{aligned}
\tag{2.10}
$$

In this formulation, when the Z_2 even superfields V_m and P are evaluated on either boundary they are the usual 4D $\mathcal{N} = 1$ SUGRA multiplet. (The real field P is the prepotential for the usual conformal compensator: $\Sigma = -\frac{1}{4}\bar{D}^2 P$.) This makes coupling 5D SUGRA to fields localized on the boundaries particularly simple. For details, see Ref. [15].

The Lagrangian for linearized 5D SUGRA is

$$
\mathcal{L}_{\text{5D SUGRA}} = \mathcal{L}_{\mathcal{N}=1} + \Delta\mathcal{L}_5,
\tag{2.11}
$$

where $\mathcal{L}_{\mathcal{N}=1}$ is the linearized $\mathcal{N} = 1$ SUGRA Lagrangian (see e.g. [14])[3]

$$
\begin{aligned}
\mathcal{L}_{\mathcal{N}=1} = M_5^3 \int d^4\theta \Big[&\tfrac{1}{8} V^m D^\alpha \bar{D}^2 D_\alpha V_m + \tfrac{1}{48}\left([D^\alpha, \bar{D}^{\dot{\alpha}}]V_{\alpha\dot{\alpha}}\right)^2 - (\partial^m V_m)^2 \\
&- \tfrac{1}{3}\Sigma^\dagger\Sigma + \tfrac{2i}{3}(\Sigma - \Sigma^\dagger)\partial^m V_m \Big],
\end{aligned}
\tag{2.12}
$$

and

$$
\begin{aligned}
\Delta\mathcal{L}_5 = -\tfrac{1}{2} M_5^3 \int d^4\theta \Big\{ &\Big[T^\dagger(\Sigma - i\partial_{\alpha\dot{\alpha}}V^{\dot{\alpha}\alpha}) + \text{h.c.} \Big] \\
&- \tfrac{1}{2}\Big[D^\alpha\Psi_\alpha + \bar{D}_{\dot{\alpha}}\Psi^{\dagger\dot{\alpha}} - \partial_5 P \Big]^2 \\
&+ \Big[\partial_5 V_{\alpha\dot{\alpha}} - (\bar{D}_{\dot{\alpha}}\Psi_\alpha - D_\alpha\Psi^\dagger_{\dot{\alpha}}) \Big]^2 \Big\}.
\end{aligned}
\tag{2.13}
$$

In this normalization, $M_{\text{P}}^2 = \pi r M_5^3$, where $M_{\text{P}} = 2 \times 10^{18}$ GeV is the 4D Planck scale.

[3] We use the conventions of Wess and Bagger [16].

The terms in the Lagrangian involving the brane-localized superfields X and Q are

$$\Delta\mathcal{L}_{\text{brane}} = \delta(x^5)\mathcal{L}_{4,\text{kin}}(Q) + \delta(x^5 - \ell)\mathcal{L}_{4,\text{kin}}(X), \tag{2.14}$$

where $\mathcal{L}_{4,\text{kin}}(\Phi)$ is the kinetic term for a 4D chiral superfield Φ coupled to 4D SUGRA:

$$\mathcal{L}_{4,\text{kin}}(\Phi) = \int d^4\theta \left[\Phi^\dagger\Phi + \tfrac{2i}{3}V^m\Phi^\dagger\overleftrightarrow{\partial}_m\Phi - \tfrac{1}{3}V^m K_{mn}V^n\Phi^\dagger\Phi + \cdots\right]. \tag{2.15}$$

Here we have absorbed the conformal compensator Σ into Φ. We have omitted terms $\mathcal{O}(V^3)$ and higher, as well as $\mathcal{O}(V^2)$ with derivatives acting on the chiral fields, since these do not contribute to the terms in Eq. (2.2). Finally, K_{mn} represents the quadratic terms in Eq. (3.5) and is given explicitly by:

$$K_{mn} := \tfrac{1}{8}\eta_{mn}D^\alpha\bar{D}^2D_\alpha + \tfrac{1}{48}\sigma_m^{\dot{\alpha}\alpha}\sigma_n^{\dot{\beta}\beta}[D_\alpha, \bar{D}_{\dot{\alpha}}][D_\beta, \bar{D}_{\dot{\beta}}] + \partial_m\partial_n \tag{2.16}$$

To define the propagator for quantum calculations, we must first fix the gauge. We require just the V_mV_n propagator, because the vertices from Eq. (2.15) involve only V_m. We now show that there exists a gauge fixing term that cancels the mixing between V_m and the other bulk superfields P, Ψ_α, T, and simultaneously reduces the V_m kinetic term to the simplest possible form $V^m\Box_5V_m$. To do this, we rewrite the quadratic terms in V as

$$\mathcal{L}_{5\text{D SUGRA}} = M_5^3\int d^4\theta \left[-V^m\Box_5V_m + \mathcal{Q}(\bar{D}_{\dot{\alpha}}V^{\dot{\alpha}\alpha}) + \cdots\right], \tag{2.17}$$

where

$$\mathcal{Q}(\chi^\alpha) = \tfrac{1}{24}(\chi^\alpha D^2\chi_\alpha + \text{h.c.}) - \tfrac{1}{4}\chi^\alpha(\bar{D}_{\dot{\alpha}}D_\alpha - \tfrac{1}{3}D_\alpha\bar{D}_{\dot{\alpha}})\bar{\chi}^{\dot{\alpha}}. \tag{2.18}$$

We then add the gauge fixing term

$$\Delta\mathcal{L}_{\rm gf} = -M_5^3 \int d^4\theta\, \mathcal{Q}(\mathcal{G}^\alpha), \tag{2.19}$$

where the gauge fixing function takes the form

$$\begin{aligned}\mathcal{G}^\alpha = \bar{D}_{\dot\alpha} V^{\dot\alpha\alpha} + \frac{\partial_5}{\Box_4}\left(\bar{D}^2\Psi^\alpha - \tfrac{6i}{5}\partial^{\dot\alpha\alpha}\bar{\Psi}_{\dot\alpha} - \tfrac{2}{5}[D^\alpha, \bar{D}^{\dot\alpha}]\bar{\Psi}_{\dot\alpha}\right)\\ + \tfrac{i}{5}\frac{\partial^{\dot\alpha\alpha}}{\Box_4}\bar{D}_{\dot\alpha}(\tfrac{3}{2}T^\dagger + \Sigma^\dagger).\end{aligned} \tag{2.20}$$

With this addition, we have

$$\mathcal{L}_{\rm 5D\,SUGRA} + \Delta\mathcal{L}_{\rm gf} = M_5^3 \int d^4\theta \left[-\tfrac{1}{2}V^m \Box_5 V_m\right] + \mathcal{L}(P, T, \Psi_\alpha). \tag{2.21}$$

Note that we do not need the ghost action, since we are not computing loops involving SUGRA self-couplings. Hence the ghosts decouple and do not contribute to the quantities under consideration.

2.3 Superpropagators on the Orbifold

The perturbative theory for the model with the superfield Lagrangian given by Eq. (2.11) can be completely formulated in terms of superfields with the help of supergraphs (see e.g. [13], [14]). Although we will not review these techniques, we will describe the relevant modifications to describe the brane-world scenario. In this brane-world scenario we have an S^1/Z_2 orbifold. Thus, it is convenient to write the Feynman rules in mixed 4D momentum space and 5D position space. Further, for supergraph calculations, we use the following abbreviations

$$\int_{1,\ldots,n} = \int d^4\theta_1 \cdots \int d^4\theta_n, \quad \delta_{12} = \delta^4(\theta_1 - \theta_2), \quad \mathcal{D}_1(p) = -\tfrac{1}{4}D_1^2(p), \tag{2.22}$$

where $D_\alpha(p)$ is the SUSY covariant derivative in momentum space. We omit the p argument when this leads to no ambiguity. We also note the following identities used for manipulating covariant derivatives and delta functions under superspace integrals:

$$\begin{aligned}
&\mathcal{D}_1\delta_{12} = \mathcal{D}_2\delta_{12},\\
&\delta_{12}(\mathcal{D}_1\bar{\mathcal{D}}_1\delta_{12}) = \delta_{12}(\bar{\mathcal{D}}_1\mathcal{D}_1\delta_{12}) = \delta_{12},\\
&\delta_{12}\Big[\mathcal{O}(D_1^n\bar{D}_1^m)\delta_{12}\Big] = 0 \text{ for } n<2 \text{ or } m<2,\\
&\delta_{12}(\mathcal{D}_1\bar{\mathcal{D}}_1\mathcal{D}_1\bar{\mathcal{D}}_1\delta_{12}) = \delta_{12}(\bar{\mathcal{D}}_1\mathcal{D}_1\bar{\mathcal{D}}_1\mathcal{D}_1\delta_{12}) = -p^2\delta_{12}.
\end{aligned} \tag{2.23}$$

The V_m propagator with one endpoint fixed on the visible brane is

$$\langle V_m(1, x^5 = 0)V_n(2, x^5 = y)\rangle = i\eta_{mn}\delta_{12}\Delta(p, y), \tag{2.24}$$

where the Green function $\Delta(p, y)$ satisfies the equation

$$2M_5^3(\partial_y^2 - p^2)\Delta(p, y) = -\delta(y). \tag{2.25}$$

Since V_m is an even field, the boundary conditions are $\partial_y\Delta = 0$ at the branes. In the domain $-\ell < y < \ell$, we have

$$\Delta(p, y) = \frac{1}{2M_5^3}\,\frac{\cosh(p(|y| - \ell))}{2p\sinh(p\ell)}, \tag{2.26}$$

where $p = +\sqrt{p^m p_m}$. The propagator from the visible brane to the visible brane is given by the limit $y \to 0$:

$$\Delta_{\text{vis,vis}}(p) = \frac{1}{2M_5^3}\,\frac{1}{2p\tanh(p\ell)}. \tag{2.27}$$

The propagator between the visible and hidden branes is given by the limit $y \to \ell$:

$$\Delta_{\text{vis,hid}}(p) = \frac{1}{2M_5^3}\frac{1}{2p\sinh(p\ell)}. \tag{2.28}$$

The chiral propagators localized on the brane are given by the standard 4D expression (see e.g. [13], [14])

$$\langle \Phi^\dagger(1)\Phi(2)\rangle = -\frac{i}{p^2}\mathcal{D}_1\bar{\mathcal{D}}_2\delta_{12}. \tag{2.29}$$

The vertices between chiral fields and V_m are read off from Eq. (2.15).

We now have all of the necessary ingredients to compute the coefficients c_1 and c_2. We neglect contributions due to the derivatives of $Q, Q^\dagger$ and $X, X^\dagger$, because we are only interested in corrections to scalar masses. Since c_1 and c_2 are gauge invariant, they can be computed using the simple gauge choice described above.

A few comments about the supergraph technique are in order. The standard procedure of supergraph calculations is reviewed in Refs. [13] and [14]. The main feature is that SUSY is manifest at every step. Another feature of the supergraph technique is the presence of SUSY covariant derivatives and Grassman δ-functions. The covariant derivative algebra is what simplifies the calculations in comparison to component formulations of SUSY theories. In an arbitrary supergraph, one can transfer all covariant derivatives onto one Grassman δ-function using integration by parts. This removes all integrals over anticommuting variables except one. Then one can transform the last integral over superspace to a standard Feynman integral over conventional momentum space. This is accomplished by using the rules given in Eq. (2.23). This procedure avoids calculating large numbers of diagrams that would appear in a component formulation of a SUSY theory. Furthermore, it automatically accounts the cancellations of conventional diagrams stipulated by $\mathcal{N} = 1$ SUSY.

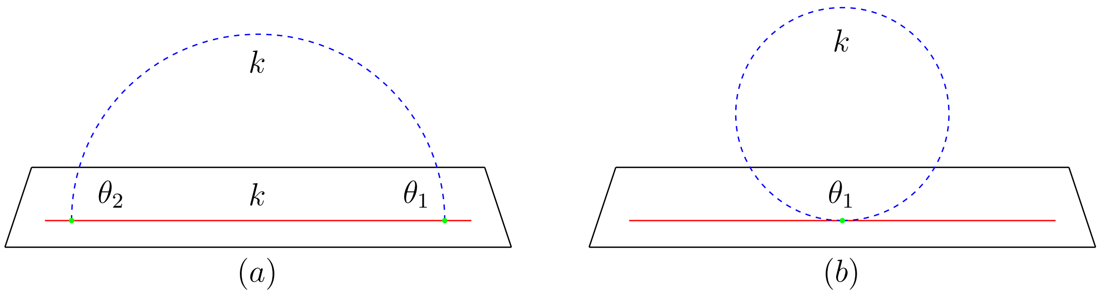

Fig. 2.1. Supergraphs contributing to the coefficient c_1, the radion mediated corrections to SUSY breaking.

2.4 Radion-mediated Contribution

We now compute the coefficient c_1 in the effective lagrangian Eq. (2.2). The 1-loop supergraphs that contribute are shown in Fig. 1.

The diagram in Fig. 1a is given by

$$\text{Fig. 1a} = -\left(\tfrac{2}{3}\right)^2 \int \frac{d^4p}{(2\pi)^4} \Delta_{\text{vis,vis}}(p) \frac{p_m p_n}{p^2}\, I^{mn}_{\text{1a}}, \tag{2.30}$$

The superspace integral I^{mn}_{1a} is

$$I^{mn}_{\text{1a}} = \eta^{mn} \int_{1,2} \delta_{12} \mathcal{D}_1 \bar{\mathcal{D}}_2 \delta_{12} = \eta^{mn} \int_1, \tag{2.31}$$

and leads to

$$\text{Fig. 1a} = -\left(\tfrac{2}{3}\right)^2 \int d^4\theta \int \frac{d^4p}{(2\pi)^4} \Delta_{\text{vis,vis}}(p). \tag{2.32}$$

The second diagram Fig. 1b gives

$$\text{Fig. 1b} = \tfrac{1}{3} \int \frac{d^4p}{(2\pi)^4} \Delta_{\text{vis,vis}}(p)\, I_{\text{1b}}. \tag{2.33}$$

For this diagram, the V_m propagator must be evaluated in the limit that θ_1 goes to θ_2. This is equivalent to inserting one more delta-function and integrating over θ_2.

Thus, the superspace integral $I_{1\mathrm{b}}$ is

$$I_{1\mathrm{b}} = \int_{1,2} \delta_{12} \eta^{mn} K_{mn} \delta_{12} = \tfrac{16}{3} \int_{1}, \tag{2.34}$$

which yields

$$\text{Fig. 1b} = \tfrac{16}{9} \int d^4\theta \int \frac{d^4 p}{(2\pi)^4} \Delta_{\mathrm{vis,vis}}(p). \tag{2.35}$$

The momentum integral is UV divergent, but its divergent part is independent of ℓ. This is easily seen from the leading behavior of the propagator at large p, which is $\Delta \to 1/(2p)$. Physically, this UV divergent contribution renormalizes the Q kinetic term on the visible brane, which is insensitive to the size of the extra dimension. For radion-mediated SUSY breaking we are interested in the ℓ dependent contribution, so we write

$$\begin{aligned}
\int \frac{d^4 p}{(2\pi)^4} \Delta_{\mathrm{vis,vis}}(p) &= \frac{1}{2M_5^3} \int \frac{d^4 p}{(2\pi)^4} \frac{1}{2p \tanh(p\ell)} \\
&= \frac{1}{2M_5^3} \int \frac{d^4 p}{(2\pi)^4} \frac{e^{-p\ell}}{2p \sinh(p\ell)} + \text{independent of } \ell \\
&= \frac{1}{8\pi^2 M_5^3} \frac{\zeta(3)}{(2\ell)^3} + \text{independent of } \ell.
\end{aligned} \tag{2.36}$$

where $\zeta(3) \simeq 1.202$ is the Riemann zeta function. Combining Eqs. (2.32), (2.35), and (2.36) the total result for c_1 is

$$c_1 = \frac{1}{6\pi^2 M_5^3} \frac{\zeta(3)}{(2\ell)^3}. \tag{2.37}$$

2.5 Brane-to-Brane Contribution

We now compute the coefficient c_2 in the effective lagrangian Eq. (2.2). The 1-loop supergraphs that contribute are shown in Fig. 2.

We first consider the diagram of Fig. 2a, consisting of four 4-point interactions. There are two possible contractions for this diagram.

One of them vanishes due to the SUSY covariant derivative algebra, and the other yields

$$\text{Fig. 2a} = \left(\tfrac{2}{3}\right)^4 \int \frac{d^4p}{(2\pi)^4} \left[\Delta_{\rm vis,hid}(p)\right]^2 I_{2\rm a}. \tag{2.38}$$

The superspace integral $I_{2\rm a}$ is

$$I_{2\rm a} = \int_{1,\ldots,4} \left(\mathcal{D}_1 \bar{\mathcal{D}}_3 \delta_{13}\right) \left(\mathcal{D}_4 \bar{\mathcal{D}}_2 \delta_{24}\right) \delta_{12}\delta_{34} = -\int_1 p^2, \tag{2.39}$$

and gives

$$\text{Fig. 2a} = -\left(\tfrac{2}{3}\right)^4 \int d^4\theta \int \frac{d^4p}{(2\pi)^4}\, p^2 \left[\Delta_{\rm vis,hid}(p)\right]^2. \tag{2.40}$$

The diagram of Fig. 2b contains two 4-point interactions

$$\text{Fig. 2b} = 2\left(\tfrac{1}{3}\right)^2 \int \frac{d^4p}{(2\pi)^4} \left[\Delta_{\rm vis,hid}(p)\right]^2 I_{2\rm b}, \tag{2.41}$$

where the superspace integral $I_{2\rm b}$ is

$$I_{2\rm b} = \int_{1,2} \delta_{12}\left(K_1^{mn} K_{nm,1} \delta_{12}\right) = \tfrac{28}{9} \int_1 p^2, \tag{2.42}$$

and leads to

$$\text{Fig. 2b} = \tfrac{56}{9}\left(\tfrac{1}{3}\right)^2 \int d^4\theta \int \frac{d^4p}{(2\pi)^4} \left[\Delta_{\rm vis,hid}(p)\right]^2 p^2. \tag{2.43}$$

The diagram Fig. 2c contains two 3-point and one 4-point interaction. There are two contractions, each giving the same contribution. We obtain

$$\text{Fig. 2c} = 2\times 2(-\tfrac{1}{3})\left(\tfrac{2}{3}\right)^2 \int \frac{d^4p}{(2\pi)^4} \frac{p_m p_n}{p^2} \left[\Delta_{\rm vis,hid}(p)\right]^2 I_{2\rm c}^{mn}, \tag{2.44}$$

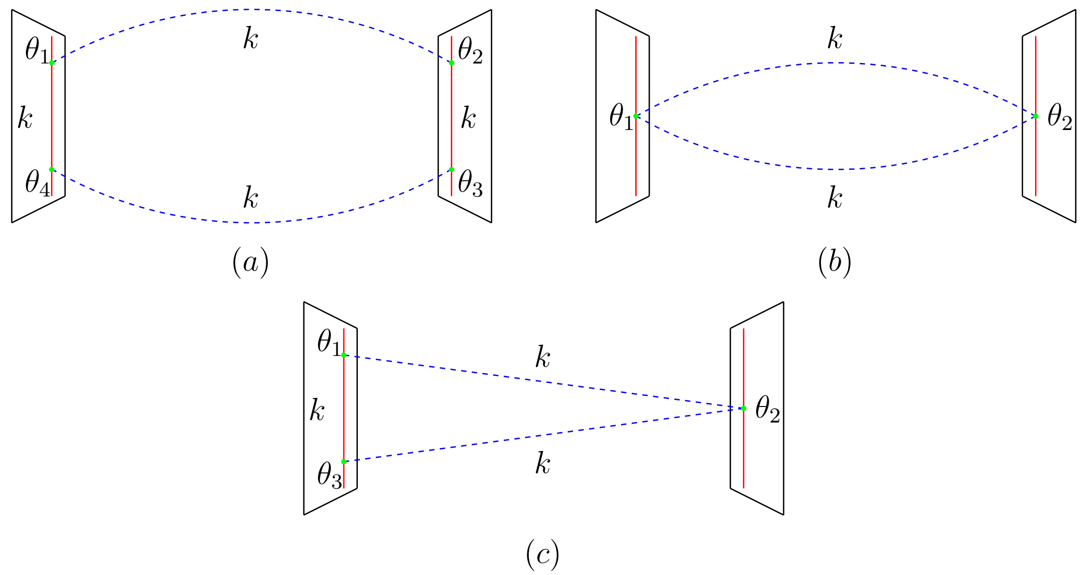

Fig. 2.2. Supergraphs contributing to the coefficient c_2, the brane-to-brane corrections to SUSY breaking.

where the superspace integral I_{2c}^{mn} is

$$I_{2\mathrm{c}}^{mn} = \int_{1,2,3} \left(\mathcal{D}_1 \bar{\mathcal{D}}_2 \delta_{12}\right) \left(K_3^{mn} \delta_{13}\right) \delta_{23} = -\tfrac{2}{3} \int_1 p^m p^n, \tag{2.45}$$

and yields

$$\text{Fig. 2c} = \tfrac{8}{9} \left(\tfrac{2}{3}\right)^2 \int \frac{d^4p}{(2\pi)^4}\, p^2 \left[\Delta_{\mathrm{vis,hid}}(p)\right]^2 . \tag{2.46}$$

The momentum integral is UV finite and can be evaluated directly. Physically, the UV finiteness is due to the fact that the SUGRA propagator cannot shrink to zero size because the endpoints are fixed on different branes. The integral is

$$\begin{aligned} \int \frac{d^4p}{(2\pi)^4}\, p^2 \left[\Delta_{\mathrm{vis,hid}}(p)\right]^2 &= \frac{1}{16 M_5^6} \int \frac{d^4p}{(2\pi)^4} \frac{1}{\sinh^2(p\ell)} \\ &= \frac{3}{16\pi^2 M_5^6} \frac{\zeta(3)}{(2\ell)^4} \end{aligned} \tag{2.47}$$

Combining Eqs. (2.40), (2.43), (2.46), and (2.47) the final result is

$$c_2 = \frac{1}{6\pi^2 M_5^6} \frac{\zeta(3)}{(2\ell)^4}. \tag{2.48}$$

This completes our calculation. The coefficients c_1 and c_2 in the 4D effective lagrangian defined in Eq. (2.2) are given by Eqs. (2.37) and (2.48), respectively.

Chapter 3: Supergravity Loop Mediated Scalar Mass Corrections

In this chapter, we investigate the supergravity loop contributions to the visible sector scalar masses from a variety of sources, either in the bulk or on the hidden brane, in a 5D 'brane-world' scenario. Extending previous work, we present a new formalism exploiting residual gauge symmetries to eliminate certain brane-localized interactions. In explicit calculations, we found a positive mass-squared contribution to the scalar masses in the case of DGP radion-mediation.

3.1 Background

In brane world scenarios, it is usual to have SUSY broken away from the visible brane and be transmitted to the visible sector by a messenger field. We shall consider the case where 5D supergravity (SUGRA) multiplet to be the sole bulk field and therefore serves as the messenger of SUSY breaking. Recent studies [19], [18] have focussed on standard brane-to-brane (where a hidden sector field breaks SUSY) and radion (where the F-term of the radion acquires a vacuum expectation value) mediations. Here we consider both these and the no-scale (where a no-scale hidden breaks SUSY) mediation. The former provides a check on a new formalism we have developed that exploits the residual gauge freedom to completely eliminate certain interactions thereby greatly simplifying the calculations. In addition, we will also explore all the above mechanisms in the context of Dvali-Gabadadze-Porratti(DGP) models and from the insights gleaned from the non-DGP case, we shall establish that both the standard brane-to-brane and no-scale mediations do not provide phenomenologically interesting results. DGP radion mediation does however provide a positive mass-squared contribution to the scalar masses.

3.2 General Setup

The Lagrangian for linearized minimal 5D SUGRA in terms of $\mathcal{N} = 1$ superfields was formulated in Ref. [15] and presented in Chapter 2. However, we repeat certain essential parts of this to point out salient features that will be useful in extending the formalism of the previous chapter.

The embedding of the 5D propagating components fields into these superfields is achieved as follows. The graviton, graviphoton and gravitino are first, respectively dimensionally reduced,

$$\begin{aligned} h_{MN} &\to h_{mn},\ h_{5m}, h_{55}, \\ B_M &\to B_m, B_5, \\ \psi_{M\tilde{\alpha}} &\to \psi^{(\pm)}_{m\alpha}, \end{aligned} \tag{3.1}$$

These reduced fields are then embedded in superfields as

$$\begin{aligned} V_m &\sim \theta\sigma^n\bar{\theta}h_{mn} + \bar{\theta}^2\theta^\alpha\psi^{(+)}_{m\alpha} + \cdots, \\ \Psi_\alpha &\sim \bar{\theta}^{\dot{\alpha}}(B_{\alpha\dot{\alpha}} + ih_{5\alpha\dot{\alpha}}) + \theta\sigma^m\bar{\theta}\psi^{(-)}_{m\alpha} + \bar{\theta}^2\psi^{(-)}_{5\alpha} + \cdots, \\ T &\sim h_{55} + iB_5 + \theta^\alpha\psi^{(+)}_{5\alpha} + \cdots. \end{aligned} \tag{3.2}$$

with the following superdiffeomorphism transformations.

$$
\begin{aligned}
\delta V_{\alpha\dot{\alpha}} &= \bar{D}_{\dot{\alpha}} L_\alpha - D_\alpha \bar{L}_{\dot{\alpha}} \\
\delta\Sigma &= -\tfrac{1}{4}\bar{D}^2 D^\alpha L_\alpha \\
\delta T &= \partial_5 \Omega \\
\delta \Psi_\alpha &= \partial_5 L_\alpha - \tfrac{1}{4} D_\alpha \Omega
\end{aligned}
\tag{3.3}
$$

In this formulation, when the Z_2 even superfields V_m and P are evaluated on either boundary they are the usual 4D $\mathcal{N}=1$ SUGRA multiplet. (The real field P is the prepotential for the usual conformal compensator: $\Sigma = -\frac{1}{4}\bar{D}^2 P$.) This makes coupling 5D SUGRA to fields localized on the boundaries particularly simple. For details, see Ref. [15].

The Lagrangian for linearized 5D SUGRA is

$$
\mathcal{L}_{\text{5D SUGRA}} = \mathcal{L}_{\mathcal{N}=1} + \Delta\mathcal{L}_5, \tag{3.4}
$$

where $\mathcal{L}_{\mathcal{N}=1}$ is the linearized $\mathcal{N}=1$ SUGRA Lagrangian (see e.g. [14])

$$
\begin{aligned}
\mathcal{L}_{\mathcal{N}=1} = M_5^3 \int d^4\theta \Big[&\tfrac{1}{8} V^m D^\alpha \bar{D}^2 D_\alpha V_m + \tfrac{1}{48}\Big([D^\alpha, \bar{D}^{\dot{\alpha}}] V_{\alpha\dot{\alpha}}\Big)^2 - (\partial^m V_m)^2 \\
&- \tfrac{1}{3}\Sigma^\dagger \Sigma + \tfrac{2i}{3}(\Sigma - \Sigma^\dagger)\partial^m V_m \Big],
\end{aligned}
\tag{3.5}
$$

and

$$
\begin{aligned}
\Delta\mathcal{L}_5 = -\tfrac{1}{2} M_5^3 \int d^4\theta \Big\{ &\Big[T^\dagger (\Sigma - i\partial_{\alpha\dot{\alpha}} V^{\dot{\alpha}\alpha}) + \text{h.c.}\Big] - \tfrac{1}{2}\Big[D^\alpha \Psi_\alpha + \bar{D}_{\dot{\alpha}} \Psi^{\dagger\dot{\alpha}} - \partial_5 P \Big]^2 \\
&+ \Big[\partial_5 V_{\alpha\dot{\alpha}} - (\bar{D}_{\dot{\alpha}} \Psi_\alpha - D_\alpha \Psi^\dagger_{\dot{\alpha}}) \Big]^2 \Big\}.
\end{aligned}
\tag{3.6}
$$

In this normalization, $M_{\rm P}^2 = \pi r M_5^3$, where $M_{\rm P} = 2 \times 10^{18}$ GeV is the 4D Planck scale.

We shall postpone further considerations of the brane-localized part of the Lagrangian till after the gauge fixing.

3.3 Gauge Fixing

We are using the gauge fixing term from [19],

$$\Delta\mathcal{L}_{\mathrm{gf}} = -M_5^3 \int d^4\theta\, \mathcal{Q}(\mathcal{G}^\alpha), \tag{3.7}$$

where the gauge fixing function takes the form

$$\mathcal{Q}(\chi^\alpha) = \tfrac{1}{24}(\chi^\alpha D^2 \chi_\alpha + \text{h.c.}) - \tfrac{1}{4}\chi^\alpha(\bar{D}_{\dot\alpha} D_\alpha - \tfrac{1}{3} D_\alpha \bar{D}_{\dot\alpha})\bar{\chi}^{\dot\alpha}. \tag{3.8}$$

and

$$\begin{aligned}\mathcal{G}^\alpha = \bar{D}_{\dot\alpha} V^{\dot\alpha\alpha} + \frac{\partial_5}{\Box_4}\left(\bar{D}^2\Psi^\alpha - \tfrac{6i}{5}\partial^{\dot\alpha\alpha}\bar{\Psi}_{\dot\alpha} - \tfrac{2}{5}[D^\alpha, \bar{D}^{\dot\alpha}]\bar{\Psi}_{\dot\alpha}\right)\\ + \tfrac{i}{5}\frac{\partial^{\dot\alpha\alpha}}{\Box_4}\bar{D}_{\dot\alpha}(\tfrac{3}{2}T^\dagger + \Sigma^\dagger)\end{aligned} \tag{3.9}$$

We shall henceforth refer to the above as the "old" gauge and the following gauge fixing as the "new" gauge.

We are interested in the residual gauge degrees of freedom so that we may eliminate certain brane-localized interactions completely. The way to proceed is to establish the bulk degrees of freedom that are fixed by the above choice of gauge-fixing. This is done by considering the variation of $\mathcal{Q}$,

$$\delta\mathcal{Q} = \tfrac{1}{12}\mathcal{G}^\alpha\left[D^2\delta\mathcal{G}_\alpha + (D_\alpha\bar{D}_{\dot\alpha} - 3\bar{D}_{\dot\alpha}D_\alpha)\delta\bar{\mathcal{G}}^{\dot\alpha}\right] + \text{c.c.} = 0 \tag{3.10}$$

After some tedious and lengthy algebra, we can simplify the above into

$$\delta\mathcal{Q} = \mathcal{G}^\alpha\left[(\Box_4 + \partial_5{}^2) - \frac{\bar{D}^2 D^2}{16\Box_4}(\Box_4 + \tfrac{4}{5}\partial_5{}^2)\right] L_\alpha + \text{c.c.} = 0 \tag{3.11}$$

where we have used the transformation properties in Eq. (3.3).

One might assume that there might be some mistake here as the rather unsightly $\frac{4}{5}$ prevents us from full factorizing the entire expression. This is the right expression for our choice of $\mathcal{Q}$. The reason we have this is because our $\mathcal{Q}$ does not diagonalize every field in the SUGRA multiplet but only V^m. The extra $\frac{1}{5}$ that is required can be thought of as coming from the gauge variation of the other fields in the SUGRA multiplet which for the purposes of this calculation is irrelevant.

$$\delta\mathcal{Q} = \mathcal{G}^\alpha \left[\frac{D^2\bar{D}^2 - 2D^\alpha \bar{D}^2 D_\alpha}{16\Box_4}(\Box_5) + \frac{\bar{D}^2 D^2}{16\Box_4}(\tfrac{1}{5}\partial_5{}^2) \right] L_\alpha + \text{c.c.} = 0 \tag{3.12}$$

Rewritten this way, we see that the linear and gauge-fixing is for the bulk and there's still brane gauge degrees of freedom are still unspecified. Let me elaborate on it. From Eq. (3.3) we see that the chiral part of L_α disappears entirely from the variation (we shall ignore $\bar{L}_{\dot\alpha}$ as it is identical to L_α except that it is the complex-conjugated analog). This means we only need consider the antichiral and linear part of L_α when we want to establish the residual gauge degrees of freedom for V^m on the brane. From Eq. (3.12), we can see that the antichiral and linear part of L_α are fully unspecified on the brane from the following.

$$\Box_5 L_{\alpha,a+l} = 0 \tag{3.13}$$

$$\Rightarrow L^{5d}_{\alpha,a+l} = L^{4d}_{\alpha,a+l} e^{\pm i\sqrt{\Box_4}x_5} \tag{3.14}$$

where $L_{\alpha,a+l} := \frac{D^2\bar{D}^2 - 2D^\alpha\bar{D}^2 D_\alpha}{16\Box_4} L_\alpha$ is the antichiral and linear part of L_α. Notice that $L^{4d}_{\alpha,a+l}$ is not specified. This means we can exploit full residual gauge freedom to eliminate certain brane-localized interactions or diagonalize DGP-type theories.

Determination of the position-momentum superpropagators is also straightforward. For the non-DGP case, i.e. both the no-scale kinetic term model as well as the standard brane-to-brane model, this follows exactly Ref.[19]. The V_m propagator is given by

$$\langle V_m(1, x^5 = 0)V_n(2, x^5 = y)\rangle = i\eta_{mn}\delta_{12}\Delta(p, y) \tag{3.15}$$

where the momentum part, for the visible brane to the visible brane, is

$$\Delta_{\text{vis,vis}}(p) = \frac{1}{2M_5^3}\frac{1}{2p\tanh(p\ell)}, \tag{3.16}$$

and for between the visible and hidden branes, is

$$\Delta_{\text{vis,hid}}(p) = \frac{1}{2M_5^3}\frac{1}{2p\sinh(p\ell)}. \tag{3.17}$$

The Q chiral propagators localized on the brane are given by the standard 4D expression:

$$\langle\Phi^\dagger(1)\Phi(2)\rangle = -\frac{i}{p^2}\mathcal{D}_1\bar{\mathcal{D}}_2\delta_{12}. \tag{3.18}$$

3.4 Standard Brane-to-brane Revisited

We revisit the calculation in Ref. [19] to test the gauge fixing in this paper and to demonstrate the utility and efficiency of this formalism.

3.4.1 Setup

To the supergravity Lagrangian we add the following terms for brane-localized superfields X and Q,

$$\Delta\mathcal{L}_{\text{brane}} = \delta(x^5)\mathcal{L}_{4,\text{kin}}(Q) + \delta(x^5 - \ell)\mathcal{L}_{4,\text{kin}}(X), \tag{3.19}$$

where $\mathcal{L}_{4,\mathrm{kin}}(\Phi)$ is the kinetic term for a 4D chiral superfield Φ coupled to 4D SUGRA:

$$\mathcal{L}_{4,\mathrm{kin}}(\Phi) = \int d^4\theta \left[\Phi^\dagger\Phi + \tfrac{2i}{3}V^m\Phi^\dagger\overleftrightarrow{\partial}_m\Phi - \tfrac{1}{3}V^m K_{mn}V^n\Phi^\dagger\Phi + \cdots\right]. \qquad (3.20)$$

where K_{mn} represents the quadratic terms in Eq. (3.5) and is given explicitly by:

$$K_{mn} := \tfrac{1}{8}\eta_{mn}D^\alpha\bar{D}^2D_\alpha + \tfrac{1}{48}\sigma_m^{\dot{\alpha}\alpha}\sigma_n^{\dot{\beta}\beta}[D_\alpha,\bar{D}_{\dot{\alpha}}][D_\beta,\bar{D}_{\dot{\beta}}] + \partial_m\partial_n \qquad (3.21)$$

3.4.2 Gauge fixing

In this particular case, we can use the original gauge-fixing from Ref. [19] by using the term

$$\Delta\mathcal{L}_{\mathrm{gf}} = -M_5^3\int d^4\theta\, \mathcal{Q}(\mathcal{G}^\alpha), \qquad (3.22)$$

where the gauge fixing function takes the form

$$\begin{aligned}\mathcal{G}^\alpha = \bar{D}_{\dot{\alpha}}V^{\dot{\alpha}\alpha} + \frac{\partial_5}{\Box_4}\left(\bar{D}^2\Psi^\alpha - \tfrac{6i}{5}\partial^{\dot{\alpha}\alpha}\bar{\Psi}_{\dot{\alpha}} - \tfrac{2}{5}[D^\alpha,\bar{D}^{\dot{\alpha}}]\bar{\Psi}_{\dot{\alpha}}\right)\\ + \tfrac{i}{5}\frac{\partial^{\dot{\alpha}\alpha}}{\Box_4}\bar{D}_{\dot{\alpha}}(\tfrac{3}{2}T^\dagger + \Sigma^\dagger).\end{aligned} \qquad (3.23)$$

This gives us the propagators defined earlier in Eq. (3.15).

Additionally, to simplify the eye diagram, we can exploit the residual gauge symmetries to replace K_{mn} by $-\eta_{mn}\Box$ on the hidden brane. This complicates the triangle as it splits it up into two parts. It does however provide a check on this method.

3.4.3 Calculation

With the above gauge fixing, we can proceed with the calculation of supersymmetry breaking contributions to scalar masses in this scenario.

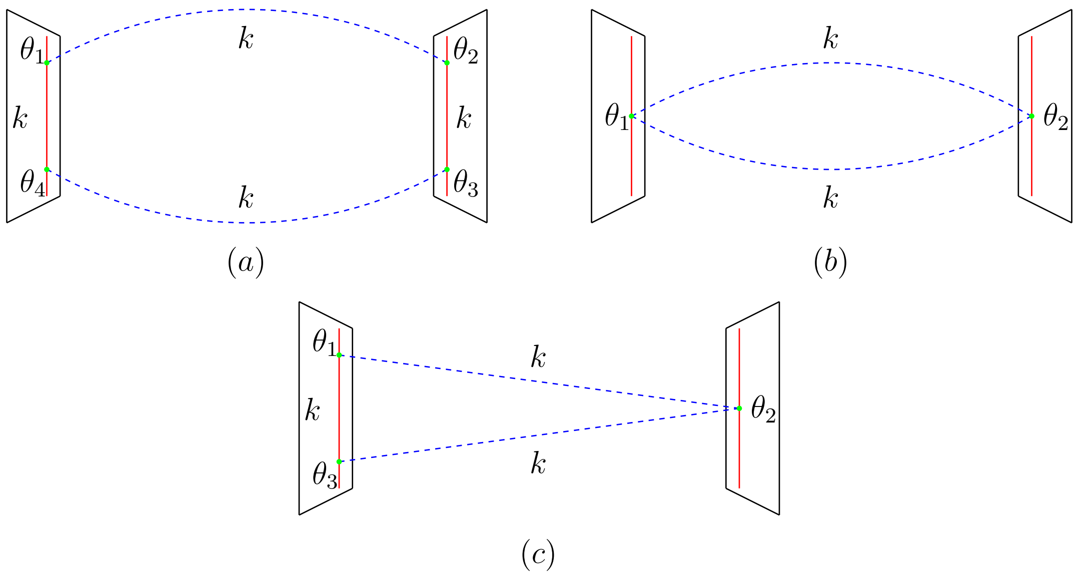

Fig. 3.1. Supergraphs contributing to the standard brane-to-brane corrections to SUSY breaking.

The bag diagram (Fig. 1a) gives us, as in Ref.[19],

$$\text{Fig. 1a} = -\left(\tfrac{2}{3}\right)^4 \int d^4\theta \int \frac{d^4p}{(2\pi)^4}\, p^2 \left[\Delta_{\text{vis,hid}}(p)\right]^2 . \tag{3.24}$$

The eye diagram (Fig. 1b) gives a different result from Ref.[19] but this is to be expected as we are working in a different gauge. The important thing, as we will show, is that the sum of all these diagrams is the same in both cases.

$$\text{Fig. 1b} = 2\left(\tfrac{1}{3}\right)^2 \int \frac{d^4p}{(2\pi)^4} \left[\Delta_{\text{vis,hid}}(p)\right]^2 p^2 I_{1\text{b}}, \tag{3.25}$$

where the superspace integral $I_{1\text{b}}$ is

$$I_{1\text{b}} = \int_{1,2} \delta_{12} \left(\text{Tr}[K^2]\delta_{12}\right) = \left(\tfrac{28}{9}, \tfrac{16}{3}\right) \int_1 p^2, \tag{3.26}$$

where the row vector $\left(\frac{28}{9}, \frac{16}{3}\right)$ represents the result in the old and new gauge respectively. We would like to point out that in the gauge adopted in Ref.[19], the above superspace integral would involve $K^{nm}K_{nm}$ which would require that we look at 9 complicated terms. In the new gauge, we replace K_{mn} by $-\eta_{mn}\Box$ on the hidden brane and the result becomes almost trivial. Putting it together,

$$\text{Fig. 1b} = \left(\frac{7.2^4}{3^4}, \frac{2^5}{3^3}\right) \int \frac{d^4p}{(2\pi)^4} \left[\Delta_{\text{vis,hid}}(p)\right]^2 p^2. \tag{3.27}$$

The triangle diagram (Fig. 1c) poses a slightly more challenging proposition. In the new gauge, the triangle diagram arising from K_{nm} localized on the visible brane (let us call this Fig. 1ci) is different from the one with K_{nm} localized on the hidden brane (Fig. 1cii) unlike Ref.[19] which in their gauge is the same. Since we are making $K_{mn} \to -\eta_{mn}\Box$ replacement on the hidden brane, this means

$$\text{Fig. 1ci} = \tfrac{4}{9}\left(\tfrac{2}{3}\right)^2 \int \frac{d^4p}{(2\pi)^4} p^2 \left[\Delta_{\text{vis,hid}}(p)\right]^2. \tag{3.28}$$

is the same as that obtained in Ref.[19].

As for Fig. 1cii,

$$\text{Fig. 1cii} = 2\tfrac{1}{3}\left(\tfrac{2}{3}\right)^2 \int \frac{d^4p}{(2\pi)^4} \left[\Delta_{\text{vis,hid}}(p)\right]^2 \frac{p^m p^n}{-p^2} I_{1\text{cii}}, \tag{3.29}$$

where the superspace integral $I_{1\text{cii}}$ is

$$I_{1\text{cii}} = \int_{1,2} \delta_{12} K_{nm} \mathcal{D}_1 \bar{\mathcal{D}}_1 \delta_{12} = \eta_{nm}\left(-\tfrac{2}{3}, 1\right) \int_1 p^2, \tag{3.30}$$

where $\mathcal{D}_1 := -\frac{D_1^2}{4}$ and $\bar{\mathcal{D}}_1 := -\frac{\bar{D}_1^2}{4}$. Once again, we have used the row vector to represent the results in the old and new gauge respectively. This yields

$$\text{Fig.1cii} = \left(\frac{2^4}{3^4}, -\frac{2^3}{3^3}\right) \int \frac{d^4p}{(2\pi)^4} \left[\Delta_{\text{vis,hid}}(p)\right]^2 p^2. \tag{3.31}$$

Combining all three diagrams,

$$\text{Fig. 1} = \tfrac{8}{9} \int \frac{d^4p}{(2\pi)^4} p^2 \left[\Delta_{\text{vis,hid}}(p)\right]^2. \tag{3.32}$$

and we see that both gauges agree as it should be.

3.5 No-scale Kinetic Term Model

In this Section, we present our calculations for the no-scale kinetic term contribution. We have actually calculated it two ways, one using the gauge fixing of Ref.[19] and the other using the gauge fixing we have developed in this paper. The latter approach is simpler, faster and emphatically demonstrates the advantages of the new formalism.

3.5.1 Setup

To the supergravity Lagrangian we add the following terms for brane-localized superfields X and Q,

$$\Delta\mathcal{L}_{\text{brane}} = \delta(x^5)\mathcal{L}_{4,\text{kin}}(Q) + \delta(x^5-\ell)\mathcal{L}_{4,\text{kin}}(X), \tag{3.33}$$

where $\mathcal{L}_{4,\text{kin}}(Q)$ is the kinetic term for a 4D chiral superfield Q coupled to 4D SUGRA,

$$\mathcal{L}_{4,\text{kin}}(Q) = \int d^4\theta \left[Q^\dagger Q + \tfrac{2i}{3}V^m Q^\dagger \overleftrightarrow{\partial}_m Q - \tfrac{1}{3}V^m K_{mn} V^n Q^\dagger Q + \cdots\right], \tag{3.34}$$

and K is defined in Eq. (3.21) and $\mathcal{L}_{4,\text{kin}}(X)$ is the no-scale kinetic term for a 4D chiral superfield X coupled to 4D SUGRA,

$$\mathcal{L}_{4,\text{kin}}(X) = \int d^4\theta \phi^\dagger\phi \left[(X^\dagger + X) - \tfrac{1}{3}V^m K_{mn} V^n (X^\dagger + X)\right] + \int d^2\theta \phi^3 W_{\text{constant}} + \text{h.c.} \tag{3.35}$$

A brief discussion is necessary here. For the case of Q, we have absorbed the conformal compensator $\phi = e^{\frac{\Sigma}{3}}$ into it. The no-scale kinetic term cannot completely absorb the conformal compensator and consequently, we should expect contact interactions of the form $V\Sigma X$. This might necessitate a more complicated gauge fixing were it not for the results from our earlier section. The other thing to note is that we have added a constant superpotential (termed a supersymmetric cosmological constant by the authors of [14]) to the Lagrangian. The reason is because we are interested in extracting the supersymmetry breaking from the F-term of the X field. In the absence of such a term, the F-term of the ϕ field would set the F-term of the X field to zero and hence be inconsistent. Such a constant superpotential is also covariant [14] but to make the term superdiffeomorphism invariant, we need to modify the transformation of Σ, [20],

$$\delta\Sigma = -\tfrac{1}{4}\bar{D}^2 D^\alpha L_\alpha - \tfrac{1}{4}\bar{D}^2(L^\alpha D_\alpha \Sigma). \tag{3.36}$$

We have also omitted higher order terms that do not contribute to the leading order of the supersymmetry breaking from the no-scale kinetic term.

3.5.2 Gauge fixing

In this particular case, we can use the old gauge from Ref. [19] or the new gauge. In the latter, one is allowed to replace the K_{mn} by $-\eta_{mn}\Box$ on one of the branes. The reason why we do not do the replacement on both branes is because brane-brane field interactions will have additional contributions through one-loop ghost diagrams. What this replacement does is that it tremendously simplifies the calculation while retaining the same form for the propagators Eq. (3.15).

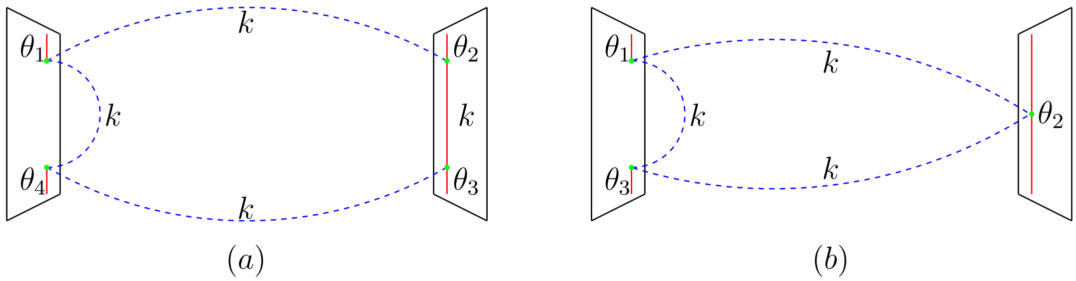

Fig. 3.2. Supergraphs contributing to the no-scale kinetic term corrections to SUSY breaking.

3.5.3 Calculation

With the above gauge fixing and determination of the propagators, we can proceed with the calculation of supersymmetry breaking contributions to scalar masses in this scenario.

The bag diagram (Fig. 2a) gives,

$$\text{Fig.2a} = -8(-\tfrac{1}{3})^2(\tfrac{2}{3})^2 \int \frac{d^4k}{(2\pi)^4}\frac{k_m k_n}{k^2}(\triangle_{\text{vis,hid}})^2(\triangle_{\text{hid,hid}})I_{2\text{a}} \tag{3.37}$$

The superspace integral is

$$I_{2\text{a}} = \int d^4\,\theta \left(\frac{\mathcal{D}\bar{\mathcal{D}}+\bar{\mathcal{D}}\mathcal{D}}{2}\right)[K^2]^{mn}\delta_{11'} = \left(\tfrac{4}{9},1\right)\eta^{mn}\,\Box_4^2 \tag{3.38}$$

where $\delta_{11'} = \lim_{2\to 1}\delta^4(\theta_1-\theta_2)$ and the row vector $\left(\frac{4}{9},1\right)$ is the result in the old and new gauge respectively.

The triangle diagram (Fig. 2b) gives,

$$\text{Fig.2b} = -8(-\tfrac{1}{3})^3 \int \frac{d^4k}{(2\pi)^4}(\triangle_{\text{vis,hid}})^2(\triangle_{\text{hid,hid}})I_{2\text{b}}. \tag{3.39}$$

The superspace integral

$$I_{2\text{b}} = \int d^4\,\theta\, \text{Tr}[K^3]\delta_{11'} = \left(\tfrac{124}{27},\tfrac{16}{3}\right)\Box_4^2 \tag{3.40}$$

where the row vector $\left(\frac{124}{27}, \frac{16}{3}\right)$ is the result in the old and new gauge respectively.

Combining these results and for both the old and new gauge fixing,

$$\text{Fig.2} = \frac{2^5}{3^3} \int \frac{d^4k}{(2\pi)^4} (\triangle_{\text{vis,hid}})^2 (\triangle_{\text{hid,hid}}) \Box_4^2 \tag{3.41}$$

we have the same result as it should be.

Wick rotate and substituting in the propagators,

$$\text{Fig.2} = \frac{2^5}{3^3} \int \frac{d^4k}{(2\pi)^4} \left(\frac{1}{2M_5^3} \frac{\cosh(k\pi R)}{2k \sinh(k\pi R)} \right)^2 \left(\frac{1}{2M_5^3} \frac{1}{2k \sinh(k\pi R)} \right) k^4 \tag{3.42}$$

Doing the momentum integral then yields,

$$\text{Fig.2} = \frac{\zeta(3)}{3^2 2^6} \frac{1}{16\pi^2 M_5^9 (\pi R)^5} \tag{3.43}$$

Alas, this is still a negative contribution to the scalar masses.

3.6 DGP-type Model

We consider the DGP-type models where there is an additional brane-localized term for the supergravity fields.

3.6.1 Setup

Consider the following DGP-type SUGRA action

$$S_{\text{DGP}} = \int d^5x \left[\mathcal{L}_{5\text{D SUGRA}} + c_v \mathcal{L}_{\mathcal{N}=1} \delta(y=0) + c_h \mathcal{L}_{\mathcal{N}=1} \delta(y=\pi R) \right] \tag{3.44}$$

where $\mathcal{L}_{5\text{D SUGRA}}$ and $\mathcal{L}_{\mathcal{N}=1}$ are defined in Eq. (3.4) and Eq. (3.5). c_h and c_v are the coefficients of the DGP terms on the hidden and visible brane respectively.

Now, depending on what we add to the branes, there are 3 interesting scenarios of supergravity loop mediated scalar mass corrections.

Case 1: Standard brane-to-brane mediation where the F-term of a hidden field (with standard canonical kinetic term).

Case 2: No-scale mediation where the F-term of a hidden field (with the no-scale canonical kinetic term) acquires a vev.

Case 3: Radion-mediation where the F-term of a radion acquires a vev.

3.6.2 Gauge fixing

From the residual gauge degrees of freedom that we have established, we know we can rewrite the DGP-type SUGRA action into a form whereby the kinetic term of V^m is always diagonal.

$$S_{\rm DGP} = -M_5^3 \int d^5x \, [V^m \Box_5 V_m + c_v V^m \Box_4 V_m \delta(y=0) + c_h V^m \Box_4 V_m \delta(y=\pi R)] . \tag{3.45}$$

This can be inverted to find the propagators. The V_m^{DGP} propagator is given by

$$\langle V_m(1, x^5 = 0) V_n(2, x^5 = y) \rangle^{DGP} = i \eta_{mn} \delta_{12} \Delta^{DGP}(p, y) \tag{3.46}$$

where the momentum part, for the visible brane to the visible brane, is

$$\Delta^{DGP}_{\rm vis,vis}(p) = \frac{1}{2M_5^3} \frac{2\, \cosh(p\pi R) + c_h\, p\, \sinh(p\pi R)}{p\, \left(2\, (c_v + c_h)\, p\, \cosh(p\pi R) + (4 + c_v\, c_h\, p^2)\, \sinh(p\pi R)\right)} \tag{3.47}$$

and for between the visible and hidden branes, is

$$\Delta^{DGP}_{\rm vis,hid}(p) = \frac{1}{2M_5^3} \frac{2}{2\, (c_v + c_h)\, p^2\, \cosh(p\pi R) + p\, (4 + c_v\, c_h\, p^2)\, \sinh(p\pi R)} . \tag{3.48}$$

Also of interest is the radius-dependent momentum part of the visible brane to the visible brane propagator,

$$\Delta^{DGP,rad}_{\text{vis,vis}}(p) = \frac{1}{2M_5^3}\,\frac{4\,(2-c_h\,p)}{e^{2p\pi R}\,p\,(2+c_v\,p)^2\,(2+c_h\,p)-p\,(c_h\,p-2)\,(c_v^2\,p^2-4)} \quad (3.49)$$

This is obtained by using method of images and instead of summing over all images, the above leaves out the "original image" as that contribution has not looped around the extra dimension and hence does not "feel" the radius. This is useful when we do the momentum integration in the case of radion mediation which requires picking up the size of the extra dimension.

Let us now consider how the DGP-type SUGRA action modifies our calculation. For Case 1, by the residual gauge freedom, the interactions (up to quadratic orders in V^m) in our model remain the same as the standard brane-to-brane non-DGP scenario. Hence, the superspace integral is identical. We can use exactly the calculation we did above and instead of plugging in the normal propagators, we use the DGP propagator, Eq. (3.48), we have just defined. Requiring that c_h and c_v be positive on both branes, we see that the momentum integral will always be positive and hence will not change the sign of the scalar mass contribution from the non-DGP scenario.

Case 2 has exactly the same interactions as the no-scale non-DGP scenario by exploiting residual gauge freedom as well. The superspace integral is unmodified and like Case 1, we plug in the DGP propagators, Eq. (3.47) and Eq. (3.48), instead of the normal propagators, Eq. (3.16) and Eq. (3.17). With c_h and c_v positive, the momentum integral has the same sign as the non-DGP scenario and hence will not give us the desired correction.

Case 3 does indeed have the same interactions as the radion-mediated non-DGP scenario. The superspace integral is unmodified, as usual, but we need to substitute in $\Delta^{DGP,rad}$ propagator which from the above form does have very interesting properties. Namely, the momentum integral can be of the opposite sign to the non-DGP case(which does not give phenomenologically interesting results). One might be concerned about the form of $\Delta^{DGP,rad}$ and say that the propagator should not flip signs as momentum is increased. However, the actual DGP propagator, $\Delta^{DGP}_{\mathrm{hid,hid}}$, is always the same sign. This includes contribution all the images (from the method of images). We know that the "original image" does not contribute to the supersymmetry breaking scalar masses and hence must be subtracted. So the resolution lies in the fact that even though the actual propagator is always physical, the part of the propagator that contributes to radion-mediation will still flip signs in certain regions of the parameter space (c_v, c_h). Additionally, we require that the propagator has a form whereby the region after the sign flip is greater than the one prior to the flip. This is so that the integral which goes into the mass correction would give us a positive effect.

3.6.3 Calculation

We shall calculate the radion-mediated DGP case which for the initial part follows that in Ref.[19].

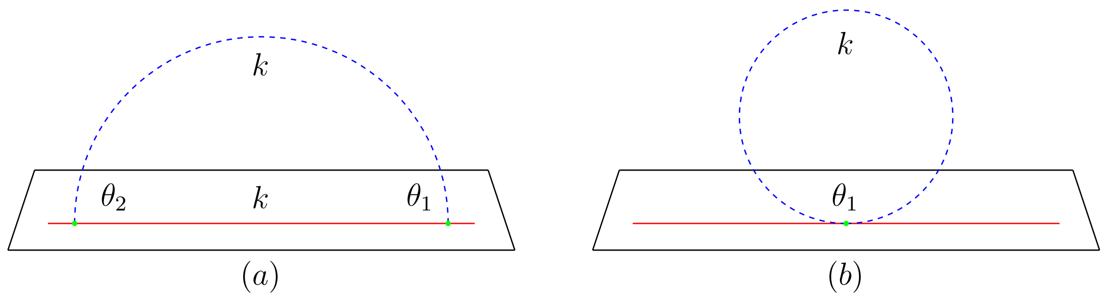

Fig. 3.3. Supergraphs contributing to the DGP radion-mediated corrections to SUSY breaking.

The diagram in Fig. 3a is given by

$$\text{Fig. 3a} = -\left(\tfrac{2}{3}\right)^2 \int \frac{d^4p}{(2\pi)^4} \Delta^{DGP,rad}_{\text{vis,vis}}(p) \frac{p_m p_n}{p^2}\, I^{mn}_{3\text{a}}, \tag{3.50}$$

The superspace integral $I^{mn}_{3\text{a}}$ is

$$I^{mn}_{3\text{a}} = \eta^{mn} \int_{1,2} \delta_{12} \mathcal{D}_1 \bar{\mathcal{D}}_2 \delta_{12} = \eta^{mn} \int_1, \tag{3.51}$$

and leads to

$$\text{Fig. 3a} = -\left(\tfrac{2}{3}\right)^2 \int \frac{d^4p}{(2\pi)^4} \Delta^{DGP,rad}_{\text{vis,vis}}(p). \tag{3.52}$$

The second diagram Fig. 1b gives

$$\text{Fig. 3b} = \tfrac{1}{3} \int \frac{d^4p}{(2\pi)^4} \Delta^{DGP,rad}_{\text{vis,vis}}(p)\, I_{3\text{b}}. \tag{3.53}$$

The superspace integral $I_{3\text{b}}$ is

$$I_{3\text{b}} = \int_{1,2} \delta_{12} \eta^{mn} K_{mn} \delta_{12} = \tfrac{16}{3} \int_1, \tag{3.54}$$

which yields

$$\text{Fig. 3b} = \tfrac{16}{9} \int \frac{d^4p}{(2\pi)^4} \Delta^{DGP,rad}_{\text{vis,vis}}(p). \tag{3.55}$$

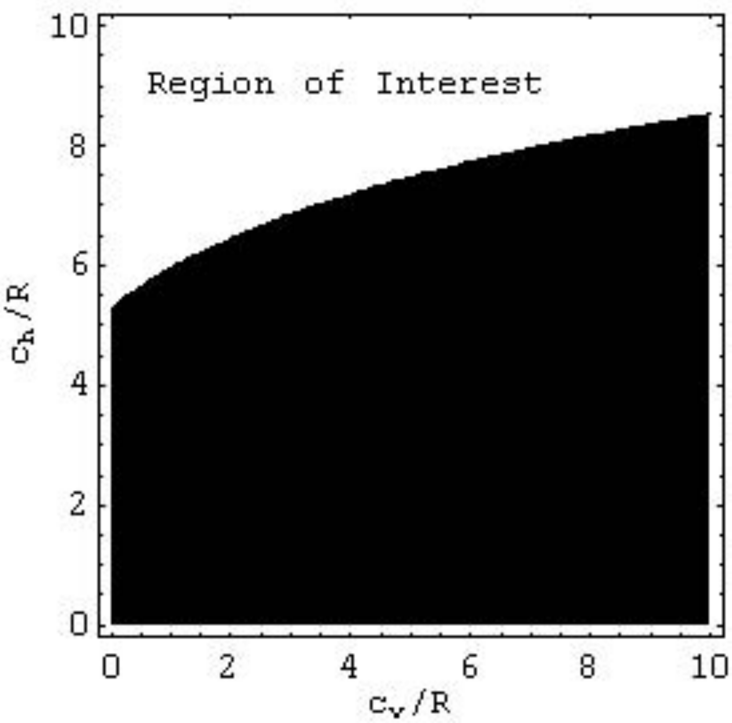

Fig. 3.4. Plot of the parameter space of the DGP radion-mediated model that gives rise to positive scalar masses. The unshaded region is the one of interest.

Since we have retained only the part of the propagator that is radius-dependent, the momentum integral is finite. The value of the integral depends on what we choose for c_v and c_h, and in Fig. 4, we have indicated the region in parameter space where one would get a positive mass contribution to the scalar masses from the DGP radion-mediated scenario.

One will notice that the positivity of the scalar masses does not depend very much on the value c_v. The reason is that the dominant contribution comes from the nearest DGP "image" (in this case the two image c_h's on the branes πR away) rather than the "original image" DGP term which has been subtracted away.

Chapter 4: Supersymmetry without Supersymmetry

We investigate the possibility that supersymmetry is not a fundamental symmetry of nature, but emerges as an accidental approximate global symmetry at low energies. This can occur if the visible sector is non-supersymmetric at high scales, but flows toward a strongly-coupled superconformal fixed point at low energies; or, alternatively, if the visible sector is localized near the infrared brane of a warped higher-dimensional spacetime with supersymmetry broken only on the UV brane. These two scenarios are related by the AdS/CFT correspondence. In order for supersymmetry to solve the hierarchy problem, the conformal symmetry must be broken below 10^{11} GeV. Accelerated unification can naturally explain the observed gauge coupling unification by physics below the conformal breaking scale. In this framework, there is no gravitino and no reason for the existence of gravitational moduli, thus eliminating the cosmological problems associated with these particles. No special dynamics is required to break supersymmetry; rather, supersymmetry is broken at observable energies because the fixed point is never reached. In 4D language, this can be due to irrelevant supersymmetry breaking operators with approximately equal dimensions. In 5D language, the size of the extra dimension is stabilized by massive bulk fields. No small input parameters are required to generate a large hierarchy. Supersymmetry can be broken in the visible sector either through direct mediation or by the F term of the modulus associated with the breaking of conformal invariance.

4.1 Background

If supersymmetry (SUSY) solves the hierarchy problem, it implies the presence of a new spacetime symmetry in nature. How does this symmetry arise? The standard paradigm is that SUSY is an exact symmetry in the fundamental UV theory, and is broken spontaneously in the IR. In this chapter, we consider the alternative possibility that UV physics is completely non-supersymmetric, and SUSY emerges as an accidental symmetry in the IR.[1]

An accidental symmetry arises when the UV theory has no relevant or marginal operators that can be added to the lagrangian to break the symmetry. In that case, all symmetry breaking effects flow to zero in the IR, and the theory becomes invariant under the symmetry at low energies even if the fundamental theory violates the symmetry maximally. A famous example is baryon and lepton number in the standard model. However, in weakly-coupled theories with scalars, scalar mass terms are always relevant, so SUSY cannot emerge as an accidental symmetry in such theories.

The situation can be very different in strongly-coupled theories. Suppose that there exists a strongly coupled superconformal theory without any relevant or marginal SUSY breaking operators that can be added to the lagrangian.[2] This

[1] The idea that 'fundamental' symmetries such as Lorentz invariance might arise as accidental symmetries of IR fixed points has been previously considered by H.B. Nielsen and collaborators [21].

[2] Every 4D CFT has a conserved stress energy operator $T_{\mu\nu}$. The traceless part of $T_{\mu\nu}$ has dimension 4, but the trace T has an anomalous dimension. T is a SUSY breaking operator that

fixed point will be attractive to all perturbations, so the boundary of the basin of attraction of the fixed point will consist of theories that have no approximate SUSY. For example, $\mathcal{N} = 1$ $SU(N)$ SUSY QCD with F flavors has a strongly-coupled fixed point near the middle of the conformal window ($F \simeq 2N$) [22]. The scalar masss operators have an uncalculable scaling dimension that is known to be larger than the canonical dimension [23]. It is possible that the anomalous dimensions are large enough that scalar masses are irrelevant in this theory, in which case this theory with the addition of large scalar mass terms flows to a superconformal fixed point in the IR.

Of course, SUSY must be broken at low energies to account for the absence of superpartners of the observed particles. The standard paradigm of exact SUSY in the UV requires special structure to break SUSY in the observable sector near the weak scale, *e.g.* dynamical SUSY breaking. While many models of dynamical SUSY breaking are known (see Ref. [24] for a review), these are very special theories and SUSY breaking is generally not robust against perturbations. The present framework offers an alternative in which the small scale of SUSY breaking in the visible sector is simply explained by the fact that the superconformal fixed point is not reached at low energies. SUSY is therefore broken explicitly rather than spontaneously; there is no Goldstino. There are several possible mechanisms that can prevent the approach to the fixed point. One possibility is that there is a relevant SUSY breaking operator with a small coefficient. This does not give an

can be added to the lagrangian, but we assume that it has a positive anomalous dimension so that it is irrelevant.

explanation of the smallness of the SUSY breaking, although the small parameter may be natural if the relevant term breaks a symmetry. In this chapter, we consider the alternative possibility that the approach to the fixed point is prevented by *irrelevant* operators. This is very natural in superconformal theories that have a moduli space of vacua in the SUSY limit. In such theories, irrelevant SUSY breaking effects will generate a potential on the moduli space, and can stabilize the moduli away from the origin. If the two lowest-dimension operators have dimensions that are somewhat close together, this can stabilize the scale modulus (dilaton) at a scale that is exponentially small compared to the fundamental scale. This mechanism is very generic, and can naturally generate a large hierarchy without small input parameters or fine tuning.

In this framework, the low-energy degrees of freedom of the visible sector are the remnants of the superconformal sector below the conformal breaking scale $\Lambda_{\rm IR}$, which is therefore the compositeness scale for the standard model matter and gauge particles. Since SUSY breaking is communicated to the visible sector only through irrelevant operators, SUSY is an approximate (but not exact) symmetry of the visible sector, *i.e.* SUSY breaking in the visible sector is naturally below the scale $\Lambda_{\rm IR}$. Direct mediation of SUSY breaking gives rise to scalar masses, gaugino masses, A terms, and μ and $B\mu$ terms, all of the same size. Understanding the absence of squark mixing requires additional structure, as in minimal supergravity. An alternative is that the CFT dynamics generates an F term for the dilaton, the modulus associated with the scale of conformal symmetry breaking. This naturally gives SUSY breaking below the scale $\Lambda_{\rm IR}$ provided that the CFT has a small parameter that breaks $U(1)_R$ symmetry. In this case, SUSY breaking in the visible sector is naturally dominated by anomaly-mediated SUSY breaking [25]. If the visible sector is the minimal supersymmetric standard model, the slepton mass-squared terms are negative. However, completely realistic non-minimal models can be constructed in this framework, using *e.g.* the ideas of Refs. [26, 27].

An important general consequence of this framework is that the gravitational sector is completely non-supersymmetric. In particular, there is no gravitino in the spectrum. This is similar to recent models in which the gravitino mass is far above the weak scale [28, 29] (see also Ref. [30]). In fact, the present framework can be thought of as a limit of the model of Ref. [28] with a high SUSY breaking scale. The absence of the gravitino eliminates the constraint on the inflationary reheat temperature that comes from the condition that gravitinos are not overproduced. Furthermore, since fundamental physics is non-supersymmetric, there there are no gravitational moduli, scalar fields with Planck-suppressed couplings that are generally present in string theory and higher-dimensional SUSY theories. This is a very good thing, because gravitational moduli cause severe cosmological difficulties that cannot be 'inflated away', and have been viewed as a major obstacle to realistic string model-building (see *e.g.* [31]). The present models do have a dilaton modulus in the 4D CFT description, but this modulus couples with more than gravitational strength, and does not give rise to cosmological difficulties.

Because gravity is not supersymmetric, gravity loops will generate SUSY breaking in the visible sector. However, these loops will be cut off at the scale Λ_{IR} where the conformal symmetry and SUSY are restored. Above the scale Λ_{IR}, the gravity loops generate a perturbation corresponding to an irrelevant operator, which is therefore suppressed by the superconformal dynamics [32]. In order for the visible scalar masses to be naturally of order 100 GeV, the scale of conformal symmetry breaking must be below 10^{11} GeV.

Because the standard model matter and gauge fields are composite below 10^{16} GeV, gauge coupling unification cannot take place in the usual way. However, accelerated unification [33] can easily lower the unification scale below the compositeness scale, thus explaining the observed unification of the standard model gauge couplings.

All of the important features of this model arise as direct consequences of strong conformal dynamics with no relevant operators. It is now well understood that 5D gravity theories in anti de Sitter (AdS) space provide 'dual' descriptions of 4D strongly-coupled conformal field theories [5, 34].[3] We can therefore write explicit 5D models that realize the framework described above. The 5D models are of the Randall–Sundrum (RS) type [35], where the UV brane breaks SUSY, while the bulk and the IR brane are supersymmetric. The visible sector is localized on the IR brane. The 5D description is weakly coupled, making explicit calculations possible. In particular, we can easily understand SUSY breaking in the visible sector. We will construct an explicit 5D model as an existence proof, but we stress that the main features are generic to superconformal theories with only irrelevant SUSY breaking operators.

[3] 5D AdS theories are invariant under $SO(4,2)$, the 4D conformal symmetry, so it is rigorously true that any 5D AdS theory describes some 4D conformal field theory.

This chapter is organized as follows. In section 2, we present an explicit 5D RS model that realizes the ideas outlined above. We consider radius stabilization and construct the low-energy 4D effective field theory, which we use to analyze SUSY breaking in the visible sector. We interpret our results in the language of 4D conformal field theories and argue that the basic features are very general. In section 3, we briefly discuss phenomenology, and section 4 contains our conclusions.

4.2 5D Model

4.2.1 Definition of the Model

Our model is based on the Randall–Sundrum (RS) model [35]. This is a 5D spacetime compactified on a S^1/Z_2 orbifold, with metric

$$ds^2 = e^{-2\sigma(y)}\eta_{\mu\nu}dx^\mu dx^\nu + dy^2, \tag{4.1}$$

where y is a periodic variable with period 2ℓ, and $\sigma(y)$ is a periodic function defined by

$$\sigma(y) = k|y| \quad \text{for} \quad -\ell < y \leq +\ell. \tag{4.2}$$

Note that this implies

$$\sigma' = \frac{d\sigma}{dy} = k\,\mathrm{sgn}(y) \quad \text{for} \quad -\ell < y \leq +\ell, \tag{4.3}$$

$$\sigma'' = \frac{d^2\sigma}{dy^2} = 2k\left[\delta(y) - \delta(y-\ell) + \cdots\right]. \tag{4.4}$$

The physical region is $0 \leq y \leq \ell$. The boundary at $y = 0$ is the 'UV brane' (or 'Planck brane') where the zero mode of graviton is localized, and the boundary at

$y = \ell$ is the 'IR brane.' We assume that the physics of this model is controlled by a single fundamental scale $\Lambda_{\rm UV} \sim M_{\rm P}$. This means that we take all couplings in the action to be of order 1 in units of $\Lambda_{\rm UV}$. The effect of the metric Eq. (4.1) is that physical scales on the IR brane are 'warped down' to the scale

$$\Lambda_{\rm IR} = \Lambda_{\rm UV} \omega, \tag{4.5}$$

where

$$\omega = e^{-k\ell} \tag{4.6}$$

is the 'warp factor.' There can be an exponentially large hierarchy between the fundamental scale and $\Lambda_{\rm IR}$, provided that the size of the extra dimension ℓ can be stabilized at a value somewhat larger than $k^{-1} \sim \Lambda_{\rm UV}^{-1}$, so that $\omega \ll 1$. This is the hierarchy generating mechanism of Randall and Sundrum [35].

The RS model is interesting in its own right, but an additional motivation to consider this model is that it can be interpreted as a strongly-coupled 4D conformal field theory (CFT) [34, 36, 37]. The origin of this equivalence is that the bulk AdS_5 geometry has a $SO(4,2)$ symmetry, which is isomorphic to the 4D conformal group, which acts on the branes as a 4D conformal transformation. In this equivalence, the bulk Kaluza–Klein (KK) modes are identified with excitations of the CFT. The UV brane acts as a UV cutoff on these modes, while the IR brane gives rise to spontaneous breaking of the conformal invariance.[4] Bulk fields are associated with

[4]More precisely, the conformal symmetry is nonlinearly realized by the position of the IR brane in the limit where the UV brane is at infinity [37].

operators of the CFT. Scalar operators that are irrelevant (respectively relevant) are associated with scalar modes with bulk mass $m^2 > 0$ (respectively $m^2 < 0$).

We are interested in 4D CFT's where the UV physics breaks SUSY, but the theory flows toward a superconformal fixed point in the IR. This means we want a 5D model where the couplings on the UV brane break SUSY maximally, while the action for the bulk and the IR brane is supersymmetric. This is radiatively stable by 5D locality. We also want the CFT to have only irrelevant perturbations. This means that all scalar fields must have positive bulk masses. We will therefore add massive hypermultiplets in the bulk, which will play an important role in stabilization and SUSY breaking.[5]

At energies below the mass of the lightest KK mode $m_{\mathrm{KK}} \sim \Lambda_{\mathrm{IR}}$ the physics can be described by a 4D effective theory that is approximately supersymmetric. The light degrees of freedom consist of the $\mathcal{N} = 1$ SUGRA multiplet, the radion chiral multiplet

$$\omega = e^{-k\ell} + \cdots + \theta^2 F_\omega, \tag{4.7}$$

and the light fields localized on the IR brane. The effective lagrangian is [38]

$$\begin{aligned}\mathcal{L}_{4,\mathrm{eff}} = -\frac{M_5^3}{k} &\int d^4\theta \, (\omega^\dagger \omega - \varphi^\dagger \varphi) \\ &+ \int d^4\theta \, \omega^\dagger \omega K_{\mathrm{IR}} + \left(\int d^2\theta \, \omega^3 W_{\mathrm{IR}} + \mathrm{h.c.} \right) \\ &+ \text{SUSY breaking terms},\end{aligned} \tag{4.8}$$

[5]The trace of the CFT stress-energy tensor corresponds to the 5D dilaton state. We therefore assume that the 5D dilaton is more massive than the hypermultiplets.

where $\varphi = 1 + \theta^2 F_\varphi$ is the conformal compensator and the superspace integrals are shorthand for the the covariant F and D projections of the superconformal tensor calculus [39]. The 4D Planck scale is

$$M_{\rm P}^2 = \frac{M_5^3}{k}(1 - \omega^2) \simeq \frac{M_5^3}{k}. \tag{4.9}$$

$K_{\rm IR}$ and $W_{\rm IR}$ are the Kähler potential and superpotential of the fields localized on the IR brane. Note that the field ω has a canonical kinetic term, and that it couples to physics on the IR brane as a dilaton. This is the nonlinear realization of the conformal symmetry, which will play an important role in what follows.

4.2.2 SUSY Breaking from 5D SUGRA

We now begin our discussion of SUSY breaking on the IR brane (the visible sector). Any SUSY breaking effects must arise by communication with the UV brane via bulk modes. In this subsection, we discuss the effects of the 5D SUGRA fields. We will discuss the effects of the bulk hypermultiplets after we have discussed stabilization.

Tree-level SUGRA KK exchange does not give rise to SUSY breaking operators on the IR brane [38, 8, 15]. There is a potential tree-level SUSY breaking effect from a constant superpotential on the IR brane. This generates a nonzero VEV for F_ω, which gives rise to anomaly-mediated SUSY breaking on the IR brane [28]. If the constant superpotential is order 1 in units of $\Lambda_{\rm UV}$, we obtain $F_\omega/\omega \sim \Lambda_{\rm IR}$, where the left-hand side is the order parameter for anomaly mediation on the IR brane. In order to obtain SUSY breaking masses at the weak scale, we must have $\Lambda_{\rm IR} \lesssim 10$ TeV. Since $\Lambda_{\rm IR}$ is also the compositeness scale for the standard model gauge fields, this implies that the standard model gauge fields are strongly coupled below 10 TeV. This requires a large number of charged states below 10 TeV, and the masses of these extra states must be finely tuned to get the observed low-energy gauge couplings. To avoid this unattractive scenario, we assume that the constant superpotential term is absent or small, which is natural by $U(1)_R$ symmetry.

We now consider SUGRA loop contributions to SUSY breaking on the IR brane, *e.g.* scalar masses. For loop momenta below $m_{\rm KK} \sim \Lambda_{\rm IR}$, the loop diagram is identical to a 4D loop diagram with a graviton line. This integral is effectively cut off for momenta above $m_{\rm KK} \sim \Lambda_{\rm IR}$ by higher-dimensional locality. This gives [28]

$$\Delta m^2_{\rm scalar} \sim \frac{1}{16\pi^2}\frac{\Lambda^4_{\rm IR}}{M^2_{\rm P}}. \tag{4.10}$$

Demanding that this contribution to the scalar masses be of order 100 GeV or less gives

$$\Lambda_{\rm IR} \lesssim 10^{11} \text{ GeV}. \tag{4.11}$$

If $\Lambda_{\rm IR} \sim 10^{11}$ GeV, this gives a flavor universal contribution to the scalar masses of order 100 GeV. In models where SUSY breaking in the visible sector is anomaly-mediated, an additional positive scalar mass-squared contribution can make the slepton masses positive in the minimal supersymmetric standard model. It would therefore be extremely interesting to compute the sign of the scalar mass contribution Eq. (4.10).[6] Gaugino masses and A terms from SUGRA loops are negligibly small.

[6]For the calculation in flat 5D theories, see Refs. [19, 18].

4.2.3 Casimir Energy

Because SUSY is broken, there will be a nonzero Casimir energy from states in the bulk that feel SUSY breaking via couplings to the UV brane. The Casimir energy depends on the radius, and therefore contributes to the radius potential.

The Casimir energy can be written as a sum over KK modes. (In the 4D CFT interpretation, these are bound states of the strongly-coupled CFT.) The radius-dependent part of the Casimir energy is UV finite, and is therefore dominated by the contribution from the lowest lying KK states, which are approximately supersymmetric. The SUSY violating mass splittings are due to gravitational strength interactions, and are therefore of order

$$\Delta m_{\rm KK} \sim \frac{m_{\rm KK}^3}{M_{\rm P}^2} \sim \omega^3. \tag{4.12}$$

The Casimir energy vanishes when the spectrum is supersymmetric, so we have

$$V_{\rm Casimir} \sim m_{\rm KK}^3 \Delta m_{\rm KK} \sim \omega^6. \tag{4.13}$$

This agrees with explicit calculations (see *e.g.* Ref. [40]).[7] We will consider stabilization mechanisms such that this contribution to the potential is negligible, so we do not need to know the sign of the Casimir energy.

4.2.4 Radius Stabilization

We now give a detailed discussion of radius stabilization in the 5D model outlined above. The 4D CFT interpretation of the radion is the scale modulus (dilaton)

[7] We thank A. Pomarol for helpful discussions on Casimir energy.

associated with spontaneous conformal symmetry breaking [37], so this is equivalent to dilaton stabilization in the CFT.

In the 5D description, the radion potential is generated by the Goldberger–Wise mechanism [41] with bulk scalar fields with positive mass-squared.[8] In the CFT description, the bulk scalars parameterize the effects of irrelevant operators in the CFT. For a single scalar, we expect a potential of the form

$$V_{\text{eff}} = a\omega^n, \tag{4.14}$$

with $n > 0$. By itself this will give a runaway potential, but we can obtain a stable minimum if there are several such terms:

$$V_{\text{eff}} = a_1\omega^{n_1} + a_2\omega^{n_2}, \tag{4.15}$$

with $n_1 > n_2 > 0$. We find a local minimum at a small value of ω if n_1 and n_2 are close in value. If we define $\epsilon = n_1 - n_2$, then

$$\omega = \left(-\frac{(n_1 - \epsilon)a_2}{n_1 a_1}\right)^{1/\epsilon} \tag{4.16}$$

is a minimum provided that $a_1 > 0$, $a_2 < 0$. (In fact, $V_{\text{eff}} < 0$ at the minimum, so this vacuum has lower than the asymptotic vacuum $\omega = 0$.) Note that ω is naturally exponentially small if the factor in parentheses is positive and less than one, and ϵ is moderately small. The potential Eq. (4.15) has the same form as the modulus potential in 'racetrack' models [42]. The radion field ω has a kinetic term

[8] Ref. [41] considered scalars with $m^2 \ll k^2$, corresponding to almost marginal 4D CFT operators; we consider $m^2 \gtrsim k^2$.

with coefficient $M_{\rm P}^2$ (see Eqs. (4.8) and (4.9)), so the physical mass of the radion is of order

$$m^2_{\rm radion} \sim \frac{\epsilon a_1}{M_{\rm P}^2}\omega^{n_1-2}. \tag{4.17}$$

We now discuss in detail how potentials of this form can arise from bulk hypermultiplets. The action for a hypermultiplet in the RS background was given in terms of $\mathcal{N}=1$ superfields by Martí and Pomarol in Ref. [43]. We add a general SUSY breaking potential on the UV brane, and a superpotential on the IR brane. These are the leading brane-localized terms in a low-energy expansion. The action is therefore

$$S_{\rm hyper} = \int d^4x \int_{-\ell}^{\ell} dy\, \mathcal{L}_5, \tag{4.18}$$

where[9]

$$\begin{aligned}\mathcal{L}_5 = \int d^4\theta\, e^{-2\sigma}(\Phi^\dagger\Phi + \tilde{\Phi}^\dagger\tilde{\Phi}) + \left[\int d^2\theta\, e^{-3\sigma}\left(\tfrac{1}{2}\tilde{\Phi}\overleftrightarrow{\partial}_y\Phi + c\sigma'\tilde{\Phi}\Phi\right) + \text{h.c.}\right] \\ - \delta(y)U(\Phi, F) + \delta(y-\ell)\,\omega^3\left[\int d^4\theta\, W(\Phi) + \text{h.c.}\right].\end{aligned} \tag{4.19}$$

The AdS/CFT correspondence relates the mass of states in the bulk to the dimension of an operator in the CFT. A 5D scalar with bulk mass m corresponds to an operator of dimension $d = 2 + \sqrt{4 + m^2/k^2}$. The bulk mass of the scalars from the hypermultiplet action above is

$$m^2_{\Phi,\tilde{\Phi}} = k^2(c \mp \tfrac{3}{2})(c \pm \tfrac{5}{2}). \tag{4.20}$$

[9]We write the action in terms of the two-sided derivative $\tilde{\Phi}\overleftrightarrow{\partial}_y\Phi = \tilde{\Phi}\partial_y\Phi - (\partial_y\tilde{\Phi})\Phi$. This is without loss of generality, since $f\tilde{\Phi}\overleftrightarrow{\partial}_y\Phi = 2f\tilde{\Phi}\partial_y\Phi + \partial_y f\tilde{\Phi}\Phi + \text{total derivative}$

so the dimensions of the operators associated with the scalar components of Φ and $\tilde{\Phi}$ are

$$\dim(\mathcal{O}_{\Phi,\tilde{\Phi}}) = 2 + |c \pm \tfrac{1}{2}|. \tag{4.21}$$

If we want the operators associated with both Φ and $\tilde{\Phi}$ to be irrelevant, we must have $|c| > \frac{5}{2}$.

For a scalar ϕ of mass m corresponding to an operator of dimension d, the general solution to the bulk equations of motion is

$$\phi = Ae^{d\sigma} + Be^{(4-d)\sigma}. \tag{4.22}$$

The coefficients A and B are fixed by the boundary conditions. The CFT interpretation of the coefficients is as follows. A is associated with a VEV of the operator

$$A = \frac{\langle \mathcal{O} \rangle}{2d - 4}, \tag{4.23}$$

while B is associated with adding a term to the CFT lagrangian

$$\Delta \mathcal{L}_{\text{CFT}} = \lambda \mathcal{O}, \tag{4.24}$$

with

$$B \propto \lambda. \tag{4.25}$$

For irrelevant operators ($d > 4$) the second term in Eq. (4.22) is exponentially decreasing in the IR, and therefore the value of the coefficient B will be determined by physics on the UV brane. In models where all dimensionful couplings are order 1 in units of Λ_{UV}, we therefore expect $B \sim \Lambda_{\text{UV}}^{3/2}$. The first term in Eq. (4.22) grows in the IR, and is therefore determined by physics on the IR brane. The scale of physics on the IR brane is set by $\Lambda_{\rm IR}$, so this will generally fix $A \sim \Lambda_{\rm IR}^{3/2} \ll \Lambda_{\rm UV}^{3/2}$. We therefore say that the first (second) term in Eq. (4.22) is IR (UV) dominated, respectively.

We now solve the equations of motion. We look for solutions depending only on y. The $\tilde{\Phi}$ equation of motion is

$$\partial_y F + (c - \tfrac{3}{2})\sigma' F = 0. \tag{4.26}$$

Imposing the condition that F is periodic and even, the most general solution is

$$F = F_0 e^{-(c-\frac{3}{2})\sigma}, \tag{4.27}$$

where F_0 is a constant of integration. We can define

$$F_{\rm UV} = F(0) = F_0, \tag{4.28}$$

$$F_{\rm IR} = F(\ell) = F_0 \omega^{c-\frac{3}{2}}. \tag{4.29}$$

The Φ equation of motion is

$$e^{-3\sigma}\partial_y \tilde{F} - (c + \tfrac{3}{2})\sigma' e^{-3\sigma}\tilde{F} = -\delta(y)\frac{\partial U}{\partial \Phi} + \delta(y-\ell)\omega^3 \frac{\partial^2 W}{\partial \Phi^2} F. \tag{4.30}$$

Imposing the condition that F is periodic and odd, the most general solution is

$$\tilde{F} = \tilde{F}_0 \frac{\sigma'}{k} e^{(c+\frac{3}{2})\sigma}, \tag{4.31}$$

where $\tilde{F}_0$ is a constant of integration. Even though $\tilde{F}$ is odd under the orbifold Z_2, it is discontinuous at the boundaries, so it effectively has a nonvanishing value on each boundary. It is convenient to define

$$\tilde{F}_{\rm UV} = \lim_{y\to 0+} \tilde{F} = \tilde{F}_0, \tag{4.32}$$

$$\tilde{F}_{\rm IR} = \lim_{y\to \ell-} \tilde{F} = \tilde{F}_0 \omega^{-(c+\frac{3}{2})}. \tag{4.33}$$

In terms of these quantities, the jump conditions at the UV and IR branes are

$$\tilde{F}_{\rm UV} = -\frac{1}{2}\frac{\partial U}{\partial \Phi_{\rm UV}}, \tag{4.34}$$

$$\tilde{F}_{\rm IR} = -\frac{1}{2}\frac{\partial^2 W}{\partial \Phi_{\rm IR}^2} F_{\rm IR}, \tag{4.35}$$

where $\Phi_{\rm UV} = \Phi(0)$, $\Phi_{\rm IR} = \Phi(\ell)$.

The $\tilde{F}$ equation of motion is

$$e^{-3\sigma}\partial_y \Phi + (c - \tfrac{3}{2})\sigma' e^{-3\sigma}\Phi + e^{-2\sigma}\tilde{F}^\dagger = 0. \tag{4.36}$$

Imposing the condition that Φ is periodic and even, the most general solution is

$$\Phi = \Phi_0 e^{-(c-\frac{3}{2})\sigma} - \frac{\tilde{F}_0^\dagger}{(2c+1)k} e^{(c+\frac{5}{2})\sigma}, \tag{4.37}$$

where Φ_0 is a constant of integration.

Finally, the F equation of motion is

$$e^{-3\sigma}\partial_y \tilde{\Phi} - (c + \tfrac{3}{2})\sigma' e^{-3\sigma}\tilde{\Phi} - e^{-2\sigma}F^\dagger = -\delta(y)\frac{\partial U}{\partial F} + \delta(y-\ell)\omega^3 \frac{\partial W}{\partial \Phi}. \tag{4.38}$$

Imposing the condition that $\tilde{\Phi}$ is periodic and odd, the most general solution is

$$\tilde{\Phi} = \frac{\sigma'}{k}\left[\tilde{\Phi}_0 e^{(c+\frac{3}{2})\sigma} - \frac{F_0^\dagger}{(2c-1)k} e^{-(c-\frac{5}{2})\sigma}\right], \tag{4.39}$$

where $\tilde{\Phi}_0$ is a constant of integration. To write the boundary conditions, we define

$$\tilde{\Phi}_{\rm UV} = \lim_{y\to 0+} \tilde{\Phi} = \tilde{\Phi}_0 - \frac{F_0^\dagger}{(2c-1)k}, \tag{4.40}$$

$$\tilde{\Phi}_{\rm IR} = \lim_{y\to \ell-} \tilde{\Phi} = \tilde{\Phi}_0 \omega^{-(c+\frac{3}{2})} - \frac{F_0^\dagger}{(2c-1)k}\omega^{c-\frac{5}{2}}. \tag{4.41}$$

The jump conditions at the UV and IR branes can then be written

$$\tilde{\Phi}_{\rm UV} = -\frac{1}{2}\frac{\partial U}{\partial F_{\rm UV}}, \tag{4.42}$$

$$\tilde{\Phi}_{\rm IR} = -\frac{1}{2}\frac{\partial W}{\partial \Phi_{\rm IR}}. \tag{4.43}$$

The fields $\tilde{\Phi}$ and $\tilde{F}$ are discontinuous at the boundaries. In our formulation with auxiliary fields, this is because the equations are first-order with delta function terms on the boundaries. Integrating out the auxiliary fields give rise to terms proportional to powers of delta functions, which naïvely are too singular to have a good continuum limit. However, supersymmetry and the orbifold projection evidently make sense out of these singular brane terms and give rise to the discontinuities at the boundaries.

To summarize, the solution is given by Eqs. (4.27), (4.31), (4.37), (4.39). The four constants of integration F_0, $\tilde{F}_0$, Φ_0, and $\tilde{\Phi}_0$ are to be determined from the four discontinuity conditions in Eqs. (4.34), (4.35), (4.42), and (4.43). Note that the jump conditions contain no explicit dependence on ω when expressed entirely in terms of UV or IR quantities.

We now use these results to write the effective potential. At the classical level, this is obtained by substituting the solutions to the equations of motion given above into the action and performing the integral over y to obtain a potential that depends on ω. Because the bulk terms are quadratic in Φ and $\tilde{\Phi}$, imposing the bulk equations reduces the effective potential to boundary terms. Using the jump conditions to simplify the result, we obtain the rather elegant expression

$$
\begin{aligned}
V_{\text{eff}} = U(\Phi_{\text{UV}},F_{\text{UV}}) &+ \left(\Phi_{\text{UV}}\tilde{F}_{\text{UV}} + \tilde{\Phi}_{\text{UV}}F_{\text{UV}} + \text{h.c.}\right) \\
&+ \omega^3\left(-\Phi_{\text{IR}}\tilde{F}_{\text{IR}} + \tilde{\Phi}_{\text{IR}}F_{\text{IR}} + \text{h.c.}\right).
\end{aligned}
\tag{4.44}
$$

There is implicit dependence on ω through the boundary values of the fields.

We now specialize to the case where $c > \frac{5}{2}$, so that the scalar components of Φ and $\tilde{\Phi}$ are associated with operators of dimension

$$
d = \dim(\mathcal{O}_\Phi) = c + \tfrac{5}{2}, \qquad \tilde{d} = \dim(\mathcal{O}_{\tilde{\Phi}}) = c + \tfrac{3}{2}, \tag{4.45}
$$

with $d, \tilde{d} > 4$. We also assume that all couplings in the lagrangian are order one in units of Λ_{UV}. Then we have

$$
F = F_{\text{UV}}e^{(4-d)\sigma}. \tag{4.46}
$$

This is UV dominated, so we parameterize it by F_{UV}. We also have

$$
\tilde{F} = \frac{\sigma'}{k}\,\tilde{F}_{\text{IR}}\omega^{\tilde{d}}e^{\tilde{d}\sigma}. \tag{4.47}
$$

This is IR dominated, so we parameterize it by $\tilde{F}_{\text{IR}}$. The solution for Φ is

$$
\Phi = \left[\Phi_{\text{UV}} + \frac{\tilde{F}^\dagger_{\text{IR}}\omega^{\tilde{d}}}{(2d-4)k}\right]e^{(4-d)\sigma} - \frac{\tilde{F}^\dagger_{\text{IR}}\omega^{\tilde{d}}}{(2d-4)k}e^{d\sigma}. \tag{4.48}
$$

We will see that $\tilde{F}_{\text{IR}} \ll \mathcal{O}(\omega)$ as a result of the jump equations, so the last term is small for all values of y. Φ is therefore UV dominated, and we parameterize it using Φ_{UV}. Finally, we have

$$
\tilde{\Phi} = \frac{\sigma'}{k}\left[\left(\tilde{\Phi}_{\text{IR}}\omega^{\tilde{d}} + \frac{F^\dagger_{\text{UV}}}{(2\tilde{d}-4)k}\omega^{2\tilde{d}-4}\right)e^{\tilde{d}\sigma} - \frac{F^\dagger_{\text{UV}}}{(2\tilde{d}-4)k}e^{(4-\tilde{d})\sigma}\right]. \tag{4.49}
$$

$\tilde{\Phi}$ is unsuppressed at both the UV and IR branes. We parameterize it by $\tilde{\Phi}_{\text{IR}}$.

The constants of integration F_{UV}, $\tilde{F}_{\mathrm{IR}}$, Φ_{UV}, and Φ_{IR} are determined by the jump equations. Eqs. (4.34) and (4.35) are

$$\tilde{F}_{\mathrm{IR}} = -\frac{1}{2} F_{\mathrm{UV}} \omega^{d-4} \frac{\partial^2 W}{\partial \Phi_{\mathrm{IR}}^2}, \tag{4.50}$$

$$\frac{\partial U}{\partial \Phi_{\mathrm{UV}}} = -2\tilde{F}_{\mathrm{IR}} \omega^{\tilde{d}}. \tag{4.51}$$

Eq. (4.50) implies that $\tilde{F}_{\mathrm{IR}} \lesssim \mathcal{O}(\omega^{d-4}) \ll \mathcal{O}(\omega)$, implying that Φ is UV dominated (see Eq. (4.48)). The jump conditions Eqs. (4.42) and (4.43) can be written

$$F_{\mathrm{UV}}^{\dagger} = \frac{(2\tilde{d}-4)k}{1-\omega^{2\tilde{d}-4}} \left[\frac{1}{2} \frac{\partial U}{\partial F_{\mathrm{UV}}} + \tilde{\Phi}_{\mathrm{IR}} \omega^{\tilde{d}} \right], \tag{4.52}$$

$$\tilde{\Phi}_{\mathrm{IR}} = -\frac{1}{2} \frac{\partial W}{\partial \Phi_{\mathrm{IR}}}. \tag{4.53}$$

From Eqs. (4.50) through (4.53) we can see that there are generically solutions with $\Phi_{\mathrm{UV}}, F_{\mathrm{UV}} = \mathcal{O}(\omega^0)$. The leading approximation for Φ_{UV} and F_{UV} is obtained by solving

$$0 = \frac{\partial U}{\partial \Phi_{\mathrm{UV}}}, \tag{4.54}$$

$$F_{\mathrm{UV}}^{\dagger} = (\tilde{d}-2)k \frac{\partial U}{\partial F_{\mathrm{UV}}}. \tag{4.55}$$

The corrections depend on the form of the IR superpotential.

Let us consider an example where

$$W = \kappa \Phi. \tag{4.56}$$

We then have

$$\tilde{\Phi}_{\mathrm{IR}} = -\tfrac{1}{2}\kappa, \qquad \tilde{F} \equiv 0, \tag{4.57}$$

while $\Phi_{\rm UV}$ and $F_{\rm UV}$ are determined by solving the equations

$$0 = \frac{\partial U}{\partial \Phi_{\rm UV}}, \tag{4.58}$$

$$F_{\rm UV}^{\dagger} = \frac{(\tilde{d}-2)k}{1-\omega^{2\tilde{d}-4}} \left(\frac{\partial U}{\partial F_{\rm UV}} - \kappa\omega^{\tilde{d}} \right). \tag{4.59}$$

Expanding in powers of ω,

$$\begin{aligned} \Phi_{\rm UV} &= \Phi_{\rm UV}^{(0)} + \omega^{\tilde{d}}\Phi_{\rm UV}^{(1)} + \mathcal{O}(\omega^{2\tilde{d}-4}), \\ F_{\rm UV} &= F_{\rm UV}^{(0)} + \omega^{\tilde{d}}F_{\rm UV}^{(1)} + \mathcal{O}(\omega^{2\tilde{d}-4}), \end{aligned} \tag{4.60}$$

we obtain the leading ω-dependent contribution to the effective potential:

$$V_{\rm eff} = \left[-\kappa F_{\rm UV}^{(0)} + \text{h.c.} \right] \omega^{\tilde{d}} + \mathcal{O}(\omega^{2\tilde{d}-4}). \tag{4.61}$$

It is clear that the coefficient of $\omega^{\tilde{d}}$ can have either sign, as required for stabilization. If we want the Casimir energy to be negligible compared to the effects discussed here, we must have $\tilde{d} < 6$.

The results above are in agreement with the expectations of the AdS/CFT correspondence. The CFT operator of lowest dimension associated with the hypermultiplet is $\mathcal{O}_{\tilde{\Phi}}$, with dimension $\tilde{d}$. If we add to the CFT lagrangian a term

$$\Delta\mathcal{L}_{\rm UV} = \tilde{\lambda}\mathcal{O}_{\tilde{\Phi}}, \tag{4.62}$$

conformal invariance implies that the effective potential has the form

$$V_{\rm eff} = \Lambda_{\rm IR}^4 f(\tilde{\lambda}_{\rm eff}(\Lambda_{\rm IR})), \tag{4.63}$$

where $\tilde{\lambda}_{\rm eff}(\mu) \sim \mu^{\tilde{d}-4}$ is the renormalized effective coupling at the scale μ. The vacuum energy vanishes for $\lambda = 0$ by SUSY, so $f(0) = 0$. Expanding in powers of

$\tilde{\lambda}$ therefore gives (using $\Lambda_{\mathrm{IR}} \sim \omega$)

$$V_{\mathrm{eff}} \sim \omega^{\tilde{d}} + \omega^{2\tilde{d}-4} + \cdots , \tag{4.64}$$

which agrees with the 5D result found above.

We see that we can obtain a potential of the form Eq. (4.15) with the addition of two hypermultiplets with mass parameter $c > \frac{5}{2}$.[10] To obtain a large hierarchy, we need the masses of the two hypermultiplets to be approximately equal, but the tuning required is only logarithmic.

4.2.5 SUSY Breaking from Stabilization

We now consider the size of SUSY breaking on the IR brane arising from the radius stabilization dynamics.

The hypermultiplet F terms can give rise to direct SUSY breaking from couplings on the IR brane of the form (in units where $\Lambda_{\mathrm{UV}} = 1$)

$$\begin{aligned} \Delta\mathcal{L}_{\mathrm{IR}} \sim & \int d^4\theta\, \Phi^\dagger\Phi \left(Q^\dagger Q + H^\dagger H\right) + \left(\int d^2\theta\, \Phi W^\alpha W_\alpha + \text{h.c.}\right) \\ & + \int d^2\theta\, \Phi Q^2 H + \text{h.c.} \\ & + \int d^4\theta \left(\Phi^\dagger H^2 + \Phi^\dagger \Phi H^2 + \text{h.c.}\right), \end{aligned} \tag{4.65}$$

where Q and H are matter and Higgs fields localized on the IR brane, and W_α is the field strength for standard model gauge fields, also localized on the IR brane. Just as in 'minimal SUGRA' models, this gives rise to scalar masses, gaugino masses, A

[10] Alternatively, we could use one one hypermultiplet with $c \simeq \frac{9}{2}$ (corresponding to $\tilde{d} \simeq 6$) together with Casimir energy.

terms, as well as μ and $B\mu$ terms naturally of the same size, namely

$$M_{\mathrm{SUSY}} \sim F_{\mathrm{IR}} \sim \omega^{d} \sim \Lambda_{\mathrm{IR}} \omega^{\tilde{d}-4}, \tag{4.66}$$

where we have restored the mass scales in the last step. Note that $M_{\mathrm{SUSY}} \ll \Lambda_{\mathrm{IR}}$, as expected. As a mechanism for SUSY breaking, this has the attraction of simplicity and elegance; in particular, it generates μ and $B\mu$ terms without additional complicated structure [44]. There is however no explanation of the absence of squark mixing required to avoid large flavor-changing neutral currents. This type of SUSY breaking therefore requires additional flavor structure at high scales, such as the models of Refs. [45].

In the 4D CFT interpretation, the effects parameterized by Eq. (4.65) represent SUSY breaking effects of the composite CFT states from irrelevant CFT operators. Like the operators in the 5D description, these are very generic effects that are expected to be present unless there are special symmetries that forbid them. It is remarkable that these effects can naturally give rise to all required SUSY breaking at the same scale.

Another potentially important source of SUSY breaking is the radion F term generated from the stabilization dynamics. The stabilization breaks SUSY, as can be seen from the nonzero value of the hypermultiplet F terms in the bulk. To compute the radion F term, we will use the technique of analytic continuation into superspace. The strategy is to write all SUSY breaking terms in the lagrangian using superfield spurions, and obtain a superfield form of the 4D effective potential. We can then find the dependence on the radion F term in the 4D effective theory by promoting ω to a chiral superfield [38].

We begin by writing the SUSY-breaking potential on the UV brane in terms of a superfield spurion:

$$\Delta\mathcal{L}_5 = \delta(y)\int d^4\theta\, S(\Phi, \Delta\Phi), \tag{4.67}$$

where

$$S(\Phi, F) = -\theta^4\, U(\Phi, F), \tag{4.68}$$

and we use the abbreviations

$$\Delta = -\tfrac{1}{4}D^2, \qquad \bar{\Delta} = -\tfrac{1}{4}\bar{D}^2. \tag{4.69}$$

The $\tilde{\Phi}$ and Φ superfield equations of motion are

$$e^{-2\sigma}\bar{\Delta}\tilde{\Phi}^\dagger + e^{-3\sigma}\left[\partial_y\Phi + (c - \tfrac{3}{2})\sigma'\Phi\right] = 0, \tag{4.70}$$

$$\begin{aligned} e^{-2\sigma}\bar{\Delta}\Phi^\dagger &+ e^{-3\sigma}\left[-\partial_y\tilde{\Phi} + (c + \tfrac{3}{2})\sigma'\tilde{\Phi}\right] \\ &= -\delta(y)\bar{\Delta}\left[\frac{\partial S}{\partial\Phi} + \Delta\left(\frac{\partial S}{\partial F}\right)\right] - \delta(y-\ell)\omega^3\frac{\partial W}{\partial\Phi}. \end{aligned} \tag{4.71}$$

The solutions are

$$\Phi = \Phi_0 e^{-(c-\frac{3}{2})\sigma} - \frac{1}{(2c+1)k}\bar{\Delta}\tilde{\Phi}_0^\dagger e^{(c+\frac{5}{2})\sigma}, \tag{4.72}$$

$$\tilde{\Phi} = \frac{\sigma'}{k}\left[\tilde{\Phi}_0 e^{(c+\frac{3}{2})\sigma} - \frac{1}{(2c-1)k}\bar{\Delta}\Phi_0^\dagger e^{-(c-\frac{5}{2})\sigma}\right], \tag{4.73}$$

where Φ_0 and $\tilde{\Phi}_0$ are superfield constants of integration, determined by the jump equations on the UV and IR branes:

$$\tilde{\Phi}_0 - \frac{1}{(2c-1)k}\bar{\Delta}\Phi_0^\dagger = \tfrac{1}{2}\bar{\Delta}\left[\frac{\partial S}{\partial\Phi} + \Delta\left(\frac{\partial S}{\partial F}\right)\right]_{\mathrm{UV}}, \tag{4.74}$$

$$\tilde{\Phi}_0\omega^{-(c-\frac{3}{2})} - \frac{1}{(2c-1)k}\bar{\Delta}\Phi_0^\dagger\omega^{c-\frac{5}{2}} = -\tfrac{1}{2}\left[\frac{\partial W}{\partial\Phi}\right]_{\mathrm{IR}}. \tag{4.75}$$

The effective potential is

$$\begin{aligned} V_{\text{eff}} = &\int d^4\theta \left[-S + \tfrac{1}{2} \left(\Phi \frac{\partial S}{\partial \Phi} + F \frac{\partial S}{\partial F} + \text{h.c.} \right) \right]_{\text{UV}} \\ &+ \int d^2\theta \, \omega^3 \left[-W + \tfrac{1}{2} \Phi \frac{\partial W}{\partial \Phi} \right]_{\text{IR}} + \text{h.c.} \end{aligned} \tag{4.76}$$

Note that this depends on ω implicitly via the solutions for Φ_0 and $\tilde{\Phi}_0$. It is easily verified that these expressions reproduce the component expressions given earlier.

The expression Eq. (4.76) can be directly analytically continued into superspace by promoting ω to a chiral superfield. To compute the VEV of F_ω, we need the linear term in F_ω in the effective potential. (The leading quadratic term in F_ω comes from Eq. (4.8).) Note that there is no contribution to the coefficient of F_ω from the first term in Eq. (4.76), which is a function of UV quantities. Even though this term depends implicitly on ω through Φ_0, all of the θ integrations must act on the explicit θ^4 in S to get a nonzero result. This immediately guarantees that $F_\omega/\omega \lesssim \Lambda_{\text{IR}}$, and agrees with the expectation that physics associated with the UV brane does not generate an F term for the radion. A similar argument shows that the Casimir energy from bulk SUGRA does not generate a coefficient for F_ω.

There is a nonzero linear term for F_ω from the second term in Eq. (4.76), which depends on IR quantities. This is

$$\begin{aligned} V_{\text{eff}} = \omega^2 \Bigg[3 &\left(-W + \tfrac{1}{2} \Phi \frac{\partial W}{\partial \Phi} \right) \\ &+ \tfrac{1}{2} \omega \frac{\partial \Phi_{\text{IR}}}{\partial \omega} \left(-\frac{\partial W}{\partial \Phi} + \Phi \frac{\partial^2 W}{\partial \Phi^2} \right) \Bigg]_{\text{IR}} F_\omega + \text{h.c.} + \cdots \end{aligned} \tag{4.77}$$

Note that the coefficient of F_ω vanishes identically if $W = \lambda \Phi^2$. This makes sense, since this term preserves conformal symmetry (λ is dimensionless).

For the example considered above, $W = \kappa\Phi$ and $\Phi_{\rm IR} \sim \omega^{d-4}$, we have

$$\frac{F_\omega}{\omega} \sim \omega\Phi_{\rm IR} \sim \Lambda_{\rm IR}\omega^{d-4}. \tag{4.78}$$

The left-hand side is the order parameter for anomaly mediated SUSY breaking (AMSB) on the IR brane [28]. Comparing to Eq. (4.66), we see that AMSB from this source is always smaller than direct SUSY breaking, even without taking into account the loop suppression factors in AMSB.

Integrating out F_ω also generates SUGRA corrections to the radion potential. This gives $\Delta V_{\rm eff} \sim \omega^{2d-4}$, which is negligible compared to the $\omega^{\tilde{d}}$ term found above.

We can obtain an additional contribution to the VEV of F_ω/ω if there is a small constant term in the superpotential on the IR brane:

$$\Delta W = C \tag{4.79}$$

gives

$$\frac{F_\omega}{\omega} \sim \Lambda_{\rm IR}\left(\frac{C}{\Lambda_{\rm UV}^3}\right). \tag{4.80}$$

As discussed above, if there are no small parameters in the theory, we have $C \sim \Lambda_{\rm UV}^3$ and the IR scale is too low. However $C \ll \Lambda_{\rm UV}^3$ is natural because C breaks a $U(1)_R$ symmetry. In the 4D CFT interpretation, the CFT has a small parameter that breaks the $U(1)_R$ symmetry and dynamically generates a small F term for the dilaton. In this scenario, SUSY breaking can be dominated by AMSB. This is attractive because it automatically gives flavor-blind SUSY breaking.

4.3 Phenomenology and Cosmology

In this section, we make some brief remarks about phenomenology and cosmology.

First we consider the radion, which is the only model-independent new degree of freedom in this framework. The radion is localized near the IR brane, and so its couplings to visible matter are suppressed by powers of $\Lambda_{\rm IR}$ rather than being Planck suppressed. The corresponding mode in 4D CFT language is the dilaton, which is a bound state of the CFT dynamics at the scale $\Lambda_{\rm IR}$. This modulus therefore couples more strongly than gravitational moduli, making the cosmology much safer. Because the radion decouples in the conformal limit, its couplings will be suppressed by loop factors, *e.g.*

$$\Delta\mathcal{L} \sim \frac{g^2}{16\pi^2}\frac{\hat{\omega}}{\Lambda_{\rm IR}} F^{\mu\nu}F_{\mu\nu}, \tag{4.81}$$

where $\hat{\omega}$ is the canonically normalized radion field and $F_{\mu\nu}$ is a standard model field strength. The radion mass is given by Eq. (4.17) with $n_1 = \tilde{d}$:

$$m_{\rm radion} \sim \Lambda_{\rm IR}\omega^{\tilde{d}-4}. \tag{4.82}$$

For the large values of $\Lambda_{\rm IR}$ we are considering, the radion effectively decouples in collider experiments.

There are two scenarios for SUSY breaking that we will discuss. We first consider the case where SUSY is broken in the visible sector by direct mediation from the CFT. In the 5D description, the SUSY breaking effects arise from the couplings in Eq. (4.65). In this case, we have (see Eq. (4.66))

$$\omega \sim (10^{-16})^{1/(\tilde{d}-3)}. \tag{4.83}$$

For example, for $\tilde{d} = 5$, we have $\omega \sim 10^{-8}$, which gives $\Lambda_{\rm IR} \sim 10^{10}$ GeV. From Eqs. (4.81) and (4.66), we see that the radion mass is of order 100 GeV, for any value of $\tilde{d}$. There is no model-independent prediction for the pattern of soft masses in this scenario. If we make the plausible assumption that the scalar masses and A terms are universal at the fundamental scale, the SUSY breaking pattern is the same as in 'minimal SUGRA.'

We now briefly consider radion cosmology in this scenario. We expect radion oscillations to dominate the universe when the temperature drops below the radion mass of order 100 GeV. The reheat temperature is of order

$$T_{\rm RH} \sim 3 \text{ TeV} \left(\frac{\Lambda_{\rm IR}}{10^{10} \text{ GeV}} \right)^{-1}. \tag{4.84}$$

This is easily large enough for a realistic cosmology.

The other scenario we discuss is that SUSY breaking is dominated by a nonzero constant superpotential. In this case, $m_{\rm radion} \gg 100$ GeV and SUSY breaking in the observable sector is anomaly-mediated. This also requires additional structure to obtain a realistic superpartner mass spectrum, but the mechanisms of *e.g.* Refs. [26, 27] can be used to obtain realistic models. The detailed phenomenology is model-dependent in this case also. Radion cosmology is easily realistic. For example, for $\tilde{d} = 5$ and $\omega = 10^{-7}$, we have $\Lambda_{\rm IR} \sim 10^{11}$ GeV, $m_{\rm radion} \sim 10$ TeV, and $T_{\rm RH} \sim 10^5$ GeV.

Chapter 5: Summary and Conclusions

In Chapters 2 and 3, we have formulated an $\mathcal{N} = 1$ supergraph approach to 5D supergravity (SUGRA) loop calculations, using the formulation of 5D SUGRA in terms of $\mathcal{N} = 1$ superfields of Ref. [15]. This formalism makes $\mathcal{N} = 1$ SUSY manifest, and makes couplings between bulk and brane fields particularly simple. In particular, there are no terms with higher powers of delta functions appearing in the calculation, as in component approaches. We have applied this formalism in Chapter 2 to compute the leading SUGRA loop contributions to visible sector scalar masses in the simplest 'brane world' scenario based on a flat 5D space compactified on a S^1/Z_2 orbifold. The terms in the effective lagrangian are defined in Eq. (2.2) and our results are given in Eqs. (2.37) and (2.48). The calculation requires the calculation of only five supergraphs. The same effective lagrangian terms have been calculated by R. Rattazzi, C.A. Scrucca, and A. Strumia using the component formulation of 5D supergravity. Our results agree [18]. In Chapter 3, we have extended the $\mathcal{N} = 1$ supergraph approach to 5D SUGRA loop calculations by exploiting the residual gauge symmetries. Applying this formalism to a variety of scenarios we find that the DGP radion mediation gives rise to positive scalar mass contributions.

There are a number of directions to extend the present results. Warped compactifications may also give positive loop contributions to scalar masses. It would also be interesting to extend the present results to higher dimensions, possibly to make direct contact with string theory, and also to construct the fully nonlinear theory. We leave these questions to future work.

In Chapter 4, we proposed a new paradigm for SUSY and SUSY breaking. In this approach, fundamental physics is completely non-supersymmetric, but the theory flows toward a supersymmetric conformal fixed point at low energies. SUSY is therefore an accidental approximate symmetry, and SUSY breaking at low energies is explicit rather than spontaneous. Remarkably, many of the features required of a realistic SUSY model follow very generically from the property that the fixed point is attractive, *i.e.* there are no relevant SUSY breaking perturbations of the fixed point. SUSY breaking at low energies naturally arises because the approach to the fixed point is halted due to a potential generated by irrelevant operators. This is contrast to the standard paradigm of spontaneous SUSY breaking which requires carefully chosen SUSY breaking sectors that are generally not robust against perturbations. Also, in the present framework all required SUSY breaking in the visible sector naturally occurs at the same scale with no small input parameters. If there is a small parameter due to an approximate broken $U(1)_R$ symmetry, SUSY breaking can be dominated by anomaly-mediated SUSY breaking.

The detailed phenomenology of visible sector SUSY breaking is model-dependent, but there are some general consequences of this framework. Because SUSY is not an exact symmetry in the UV, there is no Goldstino. Also, the gravitational sector is completely non-supersymmetric. This means that there is no gravitino, and hence no problems with gravitino cosmology. Also, we expect that there are no gravitational moduli, which cause grave cosmological difficulties in *e.g.* string theory. There is a dilaton modulus in the CFT, but it couples much stronger than gravity, and does not give rise to cosmological difficulties. Finally, this framework requires that the standard model is composite at a scale $\Lambda_{\rm IR} \lesssim 10^{11}$ GeV. This is below the unification scale, but one-step gauge coupling unification can naturally be explained by accelerated unification [33]. We conclude that this framework is an attractive alternative to the standard paradigm of spontaneously broken supersymmetry.

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